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Experimental Constraints on Higgs Boson Decays to TeV-scale Right-Handed Neutrinos

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Abstract: The existence of neutrino masses strongly suggests that right-handed neutrinos exist, but the data do not favor any particular scale for the Majorana mass parameters. Here I explore the possibility that these particles exist at the electroweak scale along with additional new physics at the TeV scale . Higher dimension operators involving right-handed neutrinos and the Higgs boson can introduce new decay modes of the Higgs boson, significantly modifying its phenomenology if it is light. With minimal flavor violation the Higgs boson cascade decays to 6 particles containing two highly displaced vertices. Each displaced vertex produces an odd number of leptons, leading to a dramatic signature of overall lepton violation at each vertex. I discuss the limits from the Tevatron, and find that they are close to exploring interesting regions of parameters, while limiting others. Moving beyond minimal flavor violation, cascade decays of the Higgs boson into as many as 14 particles can occur.

Keywords: higgs physics; beyond standard model; neutrino physics.

Contents

1. Introduction

New particles with masses of $O(m_Z)$ or less must be weakly coupled to electroweak gauge bosons and the electron vector and axial currents in order to be in agreement with electroweak precision data. Yet moderate couplings of these new states to other particles are allowed, including to the Higgs boson, or to other new states with masses of $O(\text{TeV})$. Production of the Higgs boson and new states with TeV scale masses are expected to occur at the LHC and

their decays into the lighter particles can lead to a number of striking signatures. Since the light states are weakly coupled, they may be long-lived and have macroscopic decay lengths leading to dramatic signatures of kinks in charged tracks [1] or highly displaced vertices [1] [2]. Either phenomena can occur in models of low-energy gauge mediated supersymmetry breaking [1], and displaced vertices leading to events with high multiplicity are predicted to occur in "hidden valley" models [2]. Direct decays of the Higgs boson into these light states can significantly affect limits and search strategies [3]. Supersymmetry without R-parity can also have striking signatures with displaced vertices occurring in Higgs boson decay [4].

Virtual effects of the new physics at the TeV scale may also be important, since the new physics presumably interacts (some what strongly) with the Higgs boson, and probably top quarks. Integrating out the new physics generates higher dimension operators involving the Higgs boson, top quark, and any other light particles in the low energy theory. Those operators that involve only the Standard Model particles and the Higgs boson may significantly affect electroweak LHC phenomenology and electroweak precision measurements [5][6] [7].

In [8] I described a scenario in which one important property of the Higgs boson - its lifetime - is significantly modified by the existence of higher dimension operators at the TeV scale and additional light states. In short, right-handed neutrinos are assumed to exist with masses around the electroweak scale [1]. Active neutrinos are assumed to acquire mass through a tree-level see-saw mechanism, which necessitates tiny neutrino Yukawa couplings. The interesting possibility of vanishing tree-level masses, which can occur for certain neutrino couplings and nearly degenerate right-handed neutrinos, is not explored here [9]. Under those circumstances active neutrino masses are generated at one-loop, allowing for larger neutrino couplings. Then the Higgs boson could predominately decay into a right-handed and a left-handed neutrino [9]. Here though the dominant decay mode of the Higgs boson, and the mass scales of the right-handed neutrinos and their decay channels are different from what I am exploring here.

Additional, unspecified physics at the TeV scale are assumed to interact with the Higgs boson and the right-handed neutrinos. Then below the TeV scale the new physics generates higher dimension operators involving the Higgs boson and the right-handed neutrinos. Due to these operators the Higgs boson can decay into right-handed neutrinos if they are light enough. These decays will only dominate over Standard Model decays if the Higgs boson mass is less than the WW mass threshold, and therefore only if at least one right-handed neutrino is lighter than the W gauge boson.

The dominant decay of the right-handed neutrino is into quarks and a lepton. Since the right-handed neutrinos are long-lived, two displaced vertices with average decay length any where from $O(mm - 10m)$ and larger occur, depending on the neutrino parameters of the model. In the detector this would be visible as a highly displaced vertex appearing to violate lepton number. The Higgs boson decays into 6 particles, producing two such displaced vertices. In this paper I discuss the collider bounds on such a scenario. Attention is given to

[1] For other, previous literature on electroweak scale right-handed neutrinos, see [9, 10, 11, 12]. Higher dimension operators were not considered by these authors.

limits from the Tevatron, because their kinematic reach is much larger than the LEP limit on a Standard Model Higgs boson.

A diverse set of searches for new physics are considered, because the nature of the new decay processes and the large Higgs boson production rate suggest that several searches have the potential to discover or exclude this scenario. Current Tevatron analyses are already beginning to exclude regions of parameter space having short average decay lengths ($\lesssim O(2-5cm)$) and $O(1)$ branching fractions for the Higgs boson to decay into the right-handed neutrinos. Decays of the Higgs into right-handed neutrinos having such short decay lengths probably have to be subdominant. This might occur naturally for some models, since such short decay lengths require that the right-handed neutrino be heavier than approximately m_W. If the Higgs boson is light it cannot decay into these right-handed neutrinos, although it could decay into other, lighter right-handed neutrino pairs which have longer average decay lengths. Decay lengths greater than $O(10cm)$ do not appear to be constrained, yet their production rates are close to what can currently be detected.

The outline of the paper is the following. Sections 2.1 and 2.2 describes the phenomenology of the decay of the Higgs boson and the right-handed neutrinos, reviewing some results from [8] and introducing notation. Model-independent branching fractions for decays of right-handed neutrinos into inclusive final states are derived. These are found to be quite useful in Section 3, which discusses possible Tevatron bounds on the new Higgs boson decay processes.

The phenomenology of the Higgs boson decay just described occurs when minimal flavor violation is used to determine the flavor structure of the higher dimension operators [13]. Since the flavor properties of higher dimension operators involving only right-handed neutrinos are weakly constrained, this assumption may be overly restrictive for these particles. If it is relaxed, then new Higgs boson decay channels are introduced, in which the Higgs boson can decay into as many as 14 particles. Section 4 briefly discusses this drastically different Higgs boson phenomenology.

Appendix A discusses dimension 6 operators within the context of minimal flavor violation and identifies new operators that introduce a new but rare decay mode for the Higgs boson. Appendix B provides details of the right-handed neutrino decays, with some attention given to final states having quantum interference.

2. Higgs Boson and Right-Handed Neutrino Decays

2.1 Higgs Boson Decays

At dimension 5 the only phenomenologically relevant operator involving right-handed neutrinos and the Higgs boson is

$$\frac{c_{IJ}}{\Lambda} N_I N_J H^\dagger H \ . \tag{2.1}$$

Here N_I, $I = 1, 2, 3$ are Majorana neutrinos (I will also refer to them as "right-handed neutrinos") and H is the Higgs boson with vacuum expectation value (vev) $\langle H \rangle = v/\sqrt{2} \simeq 175$

GeV. Other operators occur at dimension 5, but in the context of minimal flavor violation they are naturally and sufficiently suppressed [8].

Interestingly, these operators (2.1) can be relevant for the decay of the Higgs boson. For after electroweak symmetry breaking the Higgs boson can decay

$$h \to N_I N_J \tag{2.2}$$

with a substantial rate. With $c_{IJ} \simeq O(1)$, these decays of the Higgs boson into right-handed neutrinos can dominate over decays into bottom quarks for a wide range of scales Λ.

Whether the coefficients c_{IJ} of these operators are significant depends on the physics at the scale Λ. Experimentally they can be $O(1)$, since there are no constraints on these operators if $\Lambda \simeq O(\text{TeV})$. One reason is that these operators preserve custodial isospin and weak isospin, so there are no constraints from electroweak precision data. In fact, after electroweak symmetry breaking the only effect of these operators is to shift the masses of the right-handed neutrinos. The tiny observed left-handed neutrino masses do not constrain these operators either, since contributions to the left-handed neutrino masses obtained from inserting these operators into loops always appear with two powers of the neutrino couplings, and are therefore always subdominant[8].

Yet it is well-known that generic dimension 5 and 6 operators at the TeV scale involving only Standard Model fermions are excluded by flavor-changing and CP violating neutral current processes (FCNC), and also by the tiny values of the left-handed neutrino masses. A concern may be that any physics that suppresses such dangerous FCNC operators to acceptable levels may simultaneously suppress the operator (2.1). To investigate whether that occurs, in [8] I applied the minimal flavor violation hypothesis [13] [14] [15] to the operators (2.1). Under the assumptions of this hypothesis, dimension 6 contributions to the CP violating parameter ϵ can be suppressed with $\Lambda \simeq 5-10$ TeV [14], whereas contributions to lepton number violating processes are sufficiently suppressed for $\Lambda \simeq O(\text{ TeV})$ [15, 16]. Dimension 5 contributions to the left-handed neutrino masses are also subdominant with $\Lambda \simeq O(\text{TeV})$ [8].

"Minimal flavor violation" [13] is a hypothesis about the flavor structure of higher dimension operators and is a useful framework to adopt here. With two key assumptions, the flavor structure of higher dimension operators can be expressed in terms of the quark, lepton and neutrino Yukawa couplings, as well as the Majorana neutrino mass parameters. For the quarks and leptons, the two assumptions are that the symmetry group is maximal, and that there is only one order parameter for each flavor group. Relaxing either hypothesis can lead to dangerous flavor violation, for then the flavor structure of the higher dimension operators cannot be completely specified by the Yukawa couplings. For instance, if there is more than one spurion per symmetry group, then the linear combination determining the Yukawa couplings is in general different from the combination appearing in a higher dimension operator. Or, if the symmetry group is not maximal then there are more group invariants.

It is not surprising then that whether the operator coefficients (2.1) are suppressed or not depends on the broken flavor symmetry G_N of the right-handed neutrinos [8]. In the case

of a maximal flavor symmetry, $G_N = SU(3)_N \times U(1)'$ (see Appendix A for definitions), the minimal flavor violation hypothesis implies that

$$c_{IJ} = c\frac{[m_R]_{IJ}}{\Lambda} + \cdots \tag{2.3}$$

and is suppressed by the Majorana neutrino mass matrix m_R. The reason for the suppression is that both the Majorana masses and the operators (2.1) violate the $SU(3)_R$ symmetry, that in minimal flavor violation is assumed to be broken by only one order parameter.

If the broken flavor symmetry is smaller, then larger coefficients are allowed. For instance if instead $G_N = SO(3)$ then

$$c_{IJ} = c'\delta_{IJ} + c''\frac{[m_R]_{IJ}}{\Lambda} + \cdots \tag{2.4}$$

where now m_R is real. Here the operator coefficients are not suppressed. This too isn't surprising, since for universal couplings the operator (2.1) preserves an $SO(3)$ symmetry, and is therefore allowed to be unsuppressed. Here though the symmetry allows a large mass term for the right-handed neutrinos of $O(\Lambda)$. To avoid such a heavy right-handed neutrino I must make the technically natural assumption $m_R \ll \Lambda$.

For either scenario, the couplings c_{IJ} evaluated in the right-handed neutrino mass basis are flavor-diagonal to all orders in m_R/Λ. The decays $h \to N_I N_J$ are then flavor diagonal $(I = J)$. But depending on the broken flavor symmetry of the right-handed neutrinos, we see that the rate may either be unsuppressed and universal, or suppressed and non-universal. These distinctions affect both the branching fraction for the Higgs boson to decay into these channels, and the lepton flavor dependence of the final state. I refer the reader to [8] for additional details.

One finds that with $c_{IJ} = \delta_{IJ}$,

$$\Gamma(h \to N_I N_I) = \frac{v^2}{4\pi\Lambda^2} m_h \beta_I^3 \tag{2.5}$$

where β_I is the velocity of N_I. If instead $c_{IJ} = M_I \delta_{IJ}/\Lambda$, there is an additional suppression giving

$$\Gamma(h \to N_I N_I) = \frac{v^2}{4\pi\Lambda^2}\left(\frac{M_I}{\Lambda}\right)^2 m_h \beta_I^3 \ . \tag{2.6}$$

In Figures 1(a) and 1(b) the ratio $\mathcal{R} \equiv \sum_I \Gamma[h \to N_I N_I]/\Gamma[h \to b\bar{b}]$ is presented for $m_h = 120$ GeV, a range of M_I and scale Λ. (The dip in the plots near $M_I \simeq m_h/2$ GeV is due to phase space suppression.) To simplify the presentation of the plots, approximately universal right-handed neutrino masses $M_I \simeq M_R$ are assumed. In Figure 1(a) a universal coupling $c_{IJ} = \delta_{IJ}$ is used, whereas in Figure 1(b) a suppressed coupling $c_{IJ} = M_I \delta_{IJ}/\Lambda$ is assumed. Note that if the coupling is unsuppressed there is a large range of parameter space over which decays into right-handed neutrinos dominate decays to $b\bar{b}$. This is because the

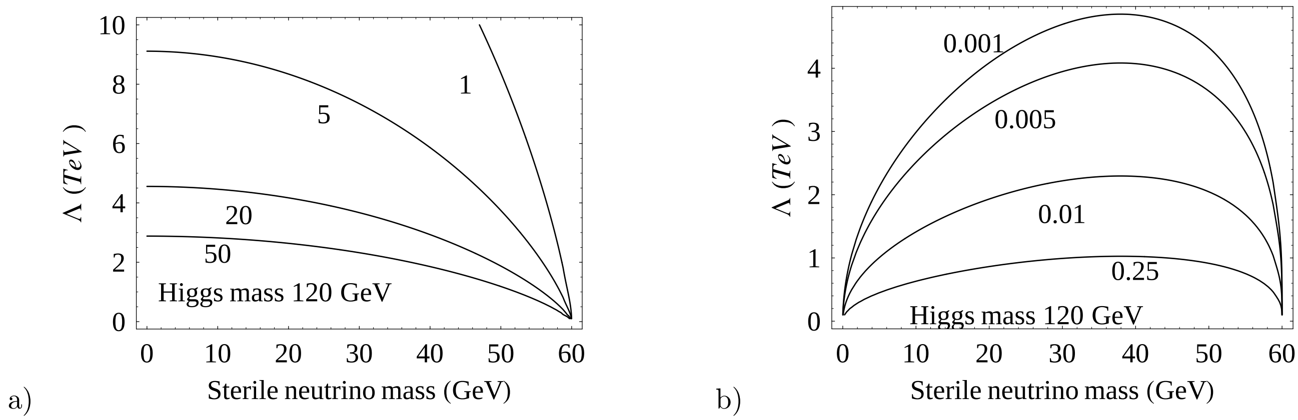

Figure 1: Contour plots of $R = \sum_I \Gamma[h \to N_I N_I]/\Gamma[h \to b\bar{b}]$ for decay into 3 flavors of right-handed neutrino, with $m_h = 120$ GeV and in (a) $c_{IJ} = \delta_{IJ}$ and (b) $c_{IJ} = M_I \delta_{IJ}/\Lambda$. Note that in (b) the scale is smaller. To simplify the presentation, right-handed neutrino masses are assumed to be universal.

decay to $b\bar{b}$ is suppressed by the small bottom Yukawa coupling. In particular, the decay to a right-handed neutrino dominates over the $b\bar{b}$ channel up to scales $\Lambda > 10$ TeV (and up to 20 TeV for $m_N \ll m_h$). (If $\mathrm{Im}(c_{IJ}) \neq 0$ the phase space for decays is larger [8]). For suppressed couplings, the ratios shown in Fig. 1(b) indicate that decays into right-handed neutrinos are subdominant, but not rare if the scale Λ is low.

If the Higgs boson mass is above both of the weak gauge boson mass thresholds, then it can decay into WW and ZZ with large rates. In this case decays of $h \to N_I N_I$ are subdominant, but are certainly still interesting. In Figures 2(a) and 2(b) the ratio $\mathcal{R} \equiv \sum_I \Gamma[h \to N_I N_I]/(\Gamma[h \to WW] + \Gamma[h \to ZZ])$ is presented for $m_h = 230$ GeV and a range of right-handed neutrino masses and scales Λ. Again, for both figures universal right-handed neutrino masses are assumed. In Figure 2(a) a universal coupling $c_{IJ} = \delta_{IJ}$ is used, whereas in Figure 2(b) a suppressed coupling $c_{IJ} = M_I \delta_{IJ}/\Lambda$ is assumed. Note that the rate for the Higgs boson to decay into all three right-handed neutrino channels is not much smaller than for a decay into the WW and ZZ channels. For example, with $\Lambda \simeq 2$ Tev and $O(1)$ coupling, the decay of the Higgs into three right-handed neutrino flavors has a branching fraction of approximately 20%, whereas for $\Lambda \simeq 10$ TeV it varies from $0.01\% - 1\%$ depending on m_R. For a suppressed coupling and a low scale, $\Lambda \simeq 2$ TeV, the branching ratio is much smaller, $O(10^{-4} - 10^{-3})$. How much integrated luminosity is needed to discover these decay processes at the LHC deserves further study.

2.2 Right-Handed Neutrino Decays

2.2.1 Dominant Decays

The standard see-saw mechanism introduces mass mixing between the right-handed and left-

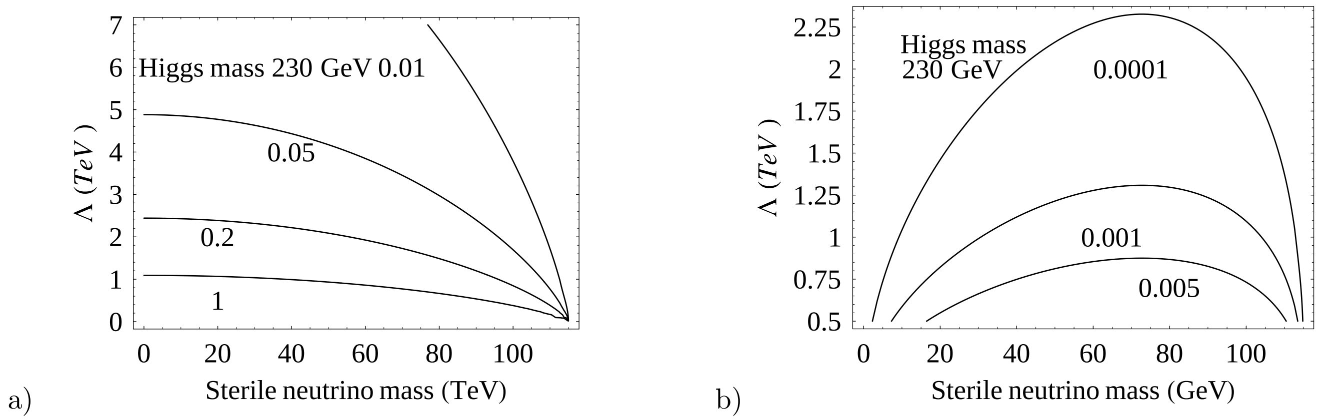

Figure 2: Contour plots of $R = \sum_I \Gamma[h \to N_I N_I]/(\Gamma[h \to WW] + \Gamma[h \to ZZ])$, with $m_h = 230$ GeV. In (a) $c_{IJ} = \delta_{IJ}$ and in (b) $c_{IJ} = \delta_{IJ} M_I/\Lambda$. Note that in (b) the scale is smaller. To simplify the presentation the right-handed neutrino masses are assumed to be degenerate.

handed neutrinos which leads to a mass matrix for the active neutrinos given by

$$m_L = \frac{1}{2}\lambda_\nu^T m_R^{-1} \lambda_\nu v^2 = m_D^T m_R^{-1} m_D \ . \tag{2.7}$$

Here λ_ν is the neutrino Yukawa coupling matrix, $m_D = \lambda_\nu v/\sqrt{2}$ is the Dirac mass matrix, and m_R is the 3×3 right-handed neutrino mass matrix. We can choose a basis where the latter matrix is diagonal and real with elements M_I. In general they will be non-universal.

The active neutrino mass matrix is diagonalized by the PMNS matrix U_{PMNS} [17] to obtain the physical masses of the active neutrinos, denoted by m_I. For generic Dirac and Majorana neutrino masses no simple relation exists between the physical masses, the rotations between the mass and gauge eigenstates, and the PMNS matrix. An approximate relation for the neutrino couplings is

$$f_I \approx 7 \times 10^{-7} \left(\frac{m_I}{0.5\text{eV}}\right)^{1/2} \left(\frac{M_I}{30\text{GeV}}\right)^{1/2} \tag{2.8}$$

where $\lambda_\nu = U_R f U_L$ is expressed in terms of two unitary matrices $U_{L/R}$ that diagonalize $\lambda_\nu^\dagger \lambda_\nu$ and $\lambda_\nu \lambda_\nu^\dagger$, and a diagonal matrix f with elements f_I. In general U_L is not the same as U_{PMNS}.

The mass mixing introduces couplings of the right-handed neutrinos to the W and Z, which at leading order in the Dirac masses are obtained by replacing a left-handed neutrino ν_I with a right-handed neutrino N_J multiplied by the mixing matrix

$$D_{IJ} = [m_D^T m_R^{-1}]_{IJ} = [m_D^T]_{IJ} M_J^{-1} \ . \tag{2.9}$$

These couplings are small, since they are suppressed by the tiny neutrino couplings. If the right-handed neutrino masses are approximately universal $M_I \simeq M$ and $U_R \simeq 1$, then

$$D_{JI} \simeq \sqrt{\frac{m_I}{M}} [U_{PMNS}]_{IJ} = 4 \times 10^{-6} \sqrt{\left(\frac{m_I}{0.5\text{eV}}\right)\left(\frac{30\text{GeV}}{M}\right)} [U_{PMNS}]_{IJ} \ . \tag{2.10}$$

Because of these couplings the right-handed neutrinos can decay into Standard Model particles, albeit with a tiny rate. If the right-handed neutrinos are heavy enough, the decays

$$N_I \to W l_J,\ Z\nu_J \tag{2.11}$$

occur. If the right-handed neutrinos are below the weak gauge boson mass threshold, the following decays into three-body final states occur,

$$N_I \to l_J W^* \to l_J f\overline{f}'\ ,\ N_I \to \nu_J Z^* \to \nu_J f\overline{f} \tag{2.12}$$

Since the right-handed neutrinos are neutral and have Majorana masses, they can also decay into the charge-conjugation of any of the final states in (2.11) and (2.12).

Not surprisingly, the dominant decay is to final states containing quarks. Model-independent branching fractions for inclusive decays are discussed below and presented in Table 2. Decays into final states containing charged leptons of specific flavor are of obvious experimental interest, but those rates are model-dependent.

Since accurate branching ratios are important for either discovery or setting limits, quantum interference effects should be included. They occur in $N_I \to l_J \bar{l}_J \nu_J$ decays, with the two leptons and the neutrino of the same flavor. This is precisely the exclusive decay studied in the D0 search for displaced muon pairs produced by the decay of a long-lived neutral particles [18]. Possible limits from their search is discussed further in Section 3.3.1. The rate for decays $N_I \to l_J \bar{l}_J \nu_J$, including the quantum interference, is discussed more in Appendix B. In the contact limit $M_I \ll m_W$ the neutral current and charged current diagrams destructively interfere, reducing the branching ratio by 0.45 compared to the incorrect result of adding the diagrams incoherently.

Quantum exchange interference also occurs in final states with three same-flavor neutrinos, $\nu_J \nu_J \overline{\nu}_J$, where the branching ratio to this final state is reduced by approximately 1/2 in the contact limit. Here the interference is less important since the rate for this final state is a fraction of the inclusive decay rate $N_I \to$ nothing. Higgs exchange interferes with neutral current exchange in decays $N_I \to b\bar{b}$+missing energy. This process has not been included and is expected to be small because of the tiny bottom Yukawa coupling.

One finds the partial decay rates

$$\Gamma[N_I \to f_J f_K f_L] + \Gamma[N_I \to f^c_J f^c_K f^c_L] = 2\frac{G_F^2 M_I^5}{192\pi^3}|D_{JI}|^2 F_{JKL} \tag{2.13}$$

Details are provided in Appendix B. The factor of "2" occurs simply because the right-handed neutrinos are Majorana particles, so at tree-level they can decay with equal rates into a state and the charge-conjugation of that state. The factor of F_{JKL} for each final state can be found in the last column of Table 1. Using the results provided in Table 1, one obtains the inclusive decay rate

$$\Gamma_{\text{total}}[N_I] = 2\frac{G_F^2 M_I^3}{192\pi^3}[m_D m_D^\dagger]_{II} N_{tot} \tag{2.14}$$

final state		F_{JKL}
$l_J q\overline{q}'$	$u\overline{d}$	$N_c c_W$
	$c\overline{s}$	$N_c c_W$
$\nu_J q\overline{q}$	$u\overline{u}+d\overline{d}+s\overline{s}$	$\sum_i \left((g_L^{(i)})^2+(g_R^{(i)})^2\right) N_c^{(i)} c_Z = 1.5c_Z$
	$c\overline{c}$	$0.43c_Z$
	$b\overline{b}$	$0.55c_Z$
$l_J\overline{l}_{K\neq J}\nu_K$		c_W
$\nu_J l_J \overline{l}_J$		$c_W + 0.13c_Z$+interference term $\rightarrow 0.59$
$\nu_J l_{K\neq J}\overline{l}_{K\neq J}$		$\left((g_L^{(K)})^2+(g_R^{(K)})^2\right) c_Z = 0.13c_Z$
$\nu_J \nu_{K\neq J}\overline{\nu}_{K\neq J}$		$\frac{1}{4}c_Z$
$\nu_J\nu_J\overline{\nu}_J$		$c_Z(\frac{1}{4}$ + quantum exchange interference) $\longrightarrow \frac{1}{8}$
total		$N_{tot}(M_I) = 8c_W + 3.2c_Z + 0.72$

Table 1: Right column gives the F factor multiplying the partial width (2.13) for the final state listed in the same row and no sum over J or K. Quantum interference between charged current and neutral current exchange occurs in final states with two charged leptons and a neutrino of the same flavor (third row from bottom). Final states with three same flavor neutrinos have a quantum exchange interference term (bottom row). In both these cases the interference is not insignificant below the gauge boson threshold. For these two final states only the F factor in the contact interaction approximation is given (denoted by the long arrows). See Appendix B for descriptions of c_W, c_Z and more details of the computations of the partial rates.

where $N_{tot}(M_I) = 8c_W + 3.2c_Z + 0.72$. The functions of c_W and c_Z are described in Appendix B and depend on the right-handed neutrino mass N_I. They are given by (B.6) and describe the effects of finite momentum transfer. (The term "0.72" in N_{tot} represents the contributions from same-flavor charged lepton $l_J\overline{l}_J\nu_J$ and same-flavor neutrino $\nu_J\nu_J\overline{\nu}_J$ final states to the total rate. Unlike all the other final states, each of these final states have quantum interfering contributions which have been computed in the low-energy limit $M_I \ll m_W$; more details can be found in Appendix B.)

Let me now comment on some of the terms appearing in the inclusive decay rate. After summing over all the final states the dependence of each partial rate on the couplings D_{JI} can be expressed in terms of the Dirac mass matrix m_D and right-handed neutrino masses M_I. This changes the overall dependence of the rate from $\Gamma \propto M_I^5$ to $\Gamma \propto M_I^3$. Additional dependence on the right-handed neutrino mass occurs through the implicit dependence of the Dirac mass on the left-handed and right-handed neutrino masses.

Inspecting Table 1 and the total decay rate (2.14) one notices that the branching fractions for some inclusive processes are independent of the Dirac mass matrix elements, since the factor of $[m_D m_D^\dagger]_{II}$ cancels in the ratio. In particular, model-independent branching fractions can be obtained for inclusive decays into quarks with any charged lepton or with missing energy, or decays into nothing. These branching fractions are listed in Table 2. The dominant decay mode is semi-leptonic, with the right-handed neutrino decaying into light

final state	Branching Fraction
light quark flavors + charged lepton	$6c_W/N_{tot} \simeq 0.50$
light quarks + missing energy	$1.5c_Z/N_{tot} \simeq 0.13$
$c\bar{c}$ + missing energy	$0.43c_Z/N_{tot} \simeq 0.036$
$b\bar{b}$ + missing energy	$0.55c_Z/N_{tot} \simeq 0.046$
two charged leptons and missing energy	$(2c_W + 0.59 + 0.26c_Z)/N_{tot} \simeq 0.24$
neutrinos	$\approx (1/8 + c_Z/2)/N_{tot} \simeq 0.05$

Table 2: Inclusive branching fractions for decays of N_I into final states containing quarks, dileptons plus missing energy, or nothing. "Charged leptons" refers to summing over e, μ and τ. These branching fractions are independent of the Dirac mass matrix elements. Processes having quantum interference were obtained in the low-energy limit $M_I \ll m_W$. Numbers displayed to the right-side of the right column are obtained in the same limit $M_I \ll m_W$, where the $c's \simeq 1$.

quark flavors and a charged lepton of any flavor, including τ. It occurs with a branching fraction of roughly 1/2. The next dominant hadronic decay mode is purely hadronic. Here the right-handed neutrino decays into light quarks and missing energy with a branching fraction of approximately 0.13. The right-handed neutrino does not have an exclusive decay into $b\bar{b}$. Instead the decay is $N_I \to b\bar{b}$+missing energy and occurs with a branching fraction of approximately 0.05. The decay $N_I \to$nothing also occurs with a branching fraction of roughly 0.05. Finally, the inclusive decay of a right-handed neutrino into missing energy and two charged leptons of opposite-sign but unspecified flavor has a branching fraction of approximately 0.24.

The rates for decays into final states containing charged leptons of specific flavors *is* model-dependent. For decays into events containing charged leptons of specific flavor J depend on the Dirac mass matrix element $|[m_D]_{JI}|^2$, which in general cannot be expressed in terms of the active neutrino masses, the right-handed neutrino masses and the PMNS matrix. This feature may be viewed as an experimental opportunity, for a measurement of these branching fractions directly determines these matrix elements, which cannot be obtained from any low-energy experiment involving the active neutrinos.

Converting the decay rate to an average lifetime gives

$$c\tau_I = 0.95m \left(\frac{1.49}{c_W + 0.40c_Z + 0.09} \right) \left(\frac{30 \text{ GeV}}{M_I} \right)^3 \left(\frac{(120 \text{ keV})^2}{[m_D m_D^\dagger]_{II}} \right) \tag{2.15}$$

This is an exact relation as no assumptions about the right-handed neutrinos masses or the Dirac matrix have been made. As already noted, diagrams having quantum interference are approximated by their low-energy values $M \ll m_W$ and represent the "0.09" contribution appearing in the formulae above [2]. The lifetime is seen to be very sensitive to the right-handed neutrino mass : changing the right-handed neutrino mass by a factor of 2 changes the average lifetime by over an order of magnitude. Longer lifetimes occur for smaller neutrino couplings, which are correlated with the active neutrino masses.

[2]As noted in [8], there final states having quantum interference were simply added incoherently. This difference accounts for the $\approx 5\%$ discrepancy between (3) and the equivalent formula given in [8].

To relate the lifetime to a measured neutrino mass requires a specific model, since the factor $[m_D m_D^\dagger]_{II} = [U_R f^2 U_R^\dagger]_{II}$ is in general not simply related to the active neutrino masses. A simple-minded approximation is $[m_D m_D^\dagger]_{II} \approx m_I M_I$ which gives

$$c\tau_I \approx 1m \left(\frac{1.49}{c_W + 0.40c_Z + 0.09}\right)\left(\frac{30 \text{ GeV}}{M}\right)^4\left(\frac{0.5\text{eV}}{m_I}\right) \tag{2.16}$$

Using (2.16), Figure 3 shows a range of active neutrino masses m_I and right-handed neutrino masses giving a variety of lifetimes.

From Figure 3 it is seen that parameters having a light Higgs boson (i.e, $m_N < m_h/2 <$ 80GeV) and $c\tau = O(5cm - 4m)$ correspond to active neutrino masses of $O(0.05 - 0.5)$eV. For most of this range the active neutrinos masses must be degenerate in order to be consistent with what is known about neutrino mass differences. This is an interesting scale, for future experiments may be sensitive to some of this region. The lower end of this range is consistent with either a normal or inverted hierarchy, with the absolute scale approximately given by the atmospheric neutrino anomaly, $\Delta m^2_{32} = (1.9 - 3.0) \times 10^{-3}\text{eV}^2 \simeq (0.05)^2\text{eV}^2$ [19].

The average decay length of a right-handed neutrino depends on its lifetime and its velocity. In the rest frame of the Higgs boson, $\gamma\beta$ for the right-handed neutrino ranges from $0.7 - 5$ for M_I/m_h ranging from $0.4 - 0.1$. Average decay lengths are then not significantly different from $O(c\tau)$ unless the right-handed neutrinos are much lighter than the Higgs boson.

I end this section by reiterating that relations between the active neutrino masses and the average decay lengths in a specific model will differ from the comments of the last paragraph and curves shown in Figure 3. The reason is that Figure 3 uses (2.16) which is a simple-minded relation between the Dirac matrix elements and the active neutrino mass. The model-independent result is given by (2.15) and depends only on $[m_D m_D^\dagger]_{II}$. While this matrix element is correlated with the active neutrino masses, the relationship is more complicated in general.

2.2.2 Subdominant Decays

It was seen that a dimension 5 operator can significantly modify the phenomenology of the Higgs boson. Could other higher dimension operators or radiative corrections cause subdominant decays that nonetheless might be important to consider? In [8] it was argued that there are no other important dimension 5 operators. At dimension 6 there are a large number of operators involving right-handed neutrinos and Standard Model particles. These may be found in Appendix A, where the size of their coefficients are estimated using minimal flavor violation. Although there are a number of operators, most of them are irrelevant either because they are suppressed by two Yukawa couplings, or because they contribute at a subdominant rate to the processes (2.11) or (2.12) previously discussed.

There is however a new process that occurs from two dimension 6 operators and a 1–loop process. After electroweak symmetry breaking there are two one–loop diagrams that together generate a magnetic and an electric moment operator for a right-handed neutrino and an active neutrino. Similarly, two of the dimension 6 operators generate the same operator,

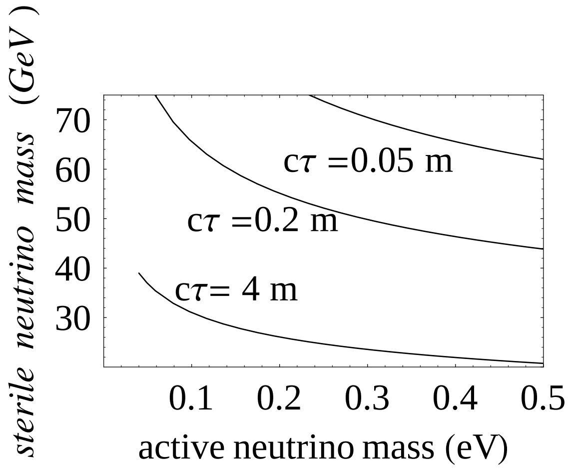

Figure 3: Average lifetime from (2.16) assuming a simple-minded relation between the Dirac matrix elements and active neutrino mass. The contours are cropped at $m_I = 0.05$ eV. For smaller values of the active neutrino mass, the right-handed neutrino mass needed to obtain $c\tau = O(0.1m - 4m)$ is off the scale and not shown for convenience.

but suppressed by the scale Λ. Together they combine into the following magnetic moment operator

$$e c_M \frac{1}{16\pi^2 v^2}(N m_D \sigma^{\mu\nu}\nu_L)F_{\mu\nu} \tag{2.17}$$

where $F_{\mu\nu}$ is the electromagnetic field strength, $\sigma^{\mu\nu} = \sigma_2(\sigma^\mu\overline{\sigma}^\nu - (\mu \leftrightarrow \nu))/4$. The coefficient of the operator is a sum of two contributions, one from a one-loop process and the other from two dimension 6 operators :

$$c_M = c_{1-loop} + c_M^{(6)}\frac{16\pi^2 v^2}{\Lambda^2} \ . \tag{2.18}$$

The one-loop contribution c_{1-loop} is calculable and expected to be $O(1)$. The coefficient $c^{(6)}$ is a linear combination of the coefficients of the dimension 6 operators described in Appendix A. The size and flavor structure of the dimension 6 operators has been estimated using minimal flavor violation. With c_{1-loop} and $c_M^{(6)}$ of the same size, c_M is dominated by the contribution from the higher dimension operators for $\Lambda \lesssim 4\pi v = 3$ TeV.

This operator causes the decay

$$N \to \gamma\nu_L \ , \ \gamma\overline{\nu}_L \ . \tag{2.19}$$

Note that these processes are not parametrically suppressed by Yukawa couplings compared to (2.11) and (2.12), since all of these amplitudes are $O(m_D)$. Summing over the three light left-handed neutrinos and anti-neutrinos, one finds the inclusive rate

$$\Gamma[N_I \to \gamma + \text{missing}] = \frac{\alpha_{em}}{2}|c_M|^2 \frac{[m_D m_D^\dagger]_{II}}{(4\pi v)^4} M_I^3 \ . \tag{2.20}$$

This rate has the same parametric dependence on the Dirac matrix elements and the right-handed neutrino mass as the rate for the dominant decay process (2.14). Comparing the two gives the following inclusive branching fraction,

$$Br(N_I \to \gamma+\text{missing}) = \alpha_{em}|c_M|^2 \frac{192\pi^3}{16}\frac{1}{(16\pi^2)^2}(c_W+0.4c_Z+0.09)^{-1} \simeq 7\times 10^{-5}|c_M|^2 \quad (2.21)$$

Like the inclusive decays given in Table 2, this branching fraction is independent of the Dirac mass matrix elements, the right-handed neutrino's mass and its flavor. Although this branching fraction may seem tiny, it is much larger at small Λ. In particular, for $\Lambda \lesssim 3$ TeV and $c_M^{(6)} \simeq c_{1-loop}$,

$$Br(N_I \to \gamma + \text{missing}) \simeq 0.007|c_M^{(6)}|^2 \left(\frac{\text{TeV}}{\Lambda}\right)^4 \quad (2.22)$$

which is sizable.

Constraints on the operator (2.17) occur from Tevatron data, which are discussed in Section 3.3. The CDF search for delayed photons [23] is beginning to exclude $c_M^{(6)} = 1$, $\Lambda \lesssim 1.3$ TeV if the dominant decay mode of the Higgs boson is to *non-relativistic* right-handed neutrinos. Because of the sensitivity of the branching fraction to the scale Λ, $\Lambda \gtrsim 2$ TeV is allowed. Constraints from low-energy experiments are weaker. For instance, although at tree-level this operator contributes to the left-handed neutrino magnetic moment because of mass-mixing between the left and right-handed neutrinos, it is tiny. One finds a neutrino magnetic moment of the size $\mu_{\nu_L} \simeq m_L m_e \mu_B/\Lambda^2 \lesssim 10^{-18}\mu_B(\text{TeV}/\Lambda)^2$ for $\Lambda \lesssim 4\pi v$, which is well-below the experimental limit (recall that in my notation m_L is the active neutrino mass matrix, not the charged lepton masses.) .

2.3 Back to Higgs Decays: Summary

Higgs bosons can decay to right-handed neutrinos with a branching fraction that depends on the Higgs boson mass, the coefficient of the operators (2.1) and the scale Λ. In the favorable situation with unsuppressed couplings and $m_h < 2m_W$, these decays can dominate over $h \to b\bar{b}$. Specific branching ratios are shown in Figures 1 and 2 for a number of scenarios. The right-handed neutrinos are long-lived due to their tiny couplings with Standard Model particles. They have displaced vertices (2.15) whose values depend on the neutrino mass parameters, and typically range from O(mm) to O(10m) or larger. The average decay length is extremely sensitive to the mass of the right-handed neutrinos 2.15, with heavier right-handed neutrinos having shorter decay lengths. The right-handed neutrinos decay predominantly into quarks and a charged lepton or missing energy. Branching fractions for some inclusive processes are given in Table 2. The right-handed neutrinos may also decay into a photon and neutrino, with an inclusive branching fraction given by (2.21). There is a calculable contribution to this decay rate which, barring accidental cancelations between it and contributions from higher dimension operators, sets a lower bound to the branching fraction for this decay.

3. Experimental Limits: Tevatron

Since the left-right mixing angles are extremely tiny, of $O(10^{-6}-10^{-5})$ and only right-handed neutrino masses $M_I > O(\text{GeV})$ are considered, constraints from neutrinoless double-β decay, measurements at the Z-pole, or from Drell-Yan production of right-handed neutrino [11], do not exist. Cosmological constraints do not exist either, since with $M_I \gtrsim O(\text{GeV})$ the right-handed neutrinos decay before Big-Bang-Nucleosynthesis.

Right-handed neutrinos are most effectively produced from the decay of a Higgs boson (see Section (2.3) for a brief summary). As discussed in the previous section, the dominant decay of the Higgs boson may be into two long-lived right-handed neutrinos, with average decay lengths $\gamma\beta c\tau \simeq O(mm - 10m)$ (or larger) in an interesting range. These decay lengths depend on the right-handed neutrino masses and the Dirac masses. Decays occur inside the detector for a wide range of parameters, but each right-handed neutrino will in general have a different average decay length.

Whether these decays dominate depends on the sizes of the coefficients of the operator (2.1). Each right-handed neutrino decays into three particles, some of which may be charged. Its dominant decay is into a charged lepton and a pair of quarks (see Table 2). The main signature for these events is a displaced quark pair and a charged lepton which reconstruct to a secondary vertex. Subdominant decays can produce two charged leptons at each secondary vertex. Further, each Higgs decay has two such secondary vertices. Since most of the decays are into visible energy, and the mass of the Higgs boson is spread across six final state particles, the event is additionally characterized by little missing energy and low p_T for each particle. Events with higher p_T and missing energy occur when the Higgs boson is produced in association with a W or Z.

The Tevatron currently does not have any direct searches for this type of signal. A direct search for this signal will probably require several complementary searches, since a wide range of average decay lengths is expected and no one search can cover all possible values. Moreover, in a given model the three right-hand neutrinos will in general have different average decay lengths, so a single search may only be sensitive to part of the spectrum.

There are a number of existing searches for new physics that have similar, although not identical, signatures. These are the searches for $R-$parity violation in supersymmetric models, searches for long-lived neutral particles, inclusive searches for events containing same-signed dileptons and searches for anomalous events containing photons. These searches are beginning to place some bounds on the parameters of the model discussed here. Given that, more accurate analyzes of the processes described below are needed. It is clear that interesting bounds could be obtained with more integrated luminosity combined with a search optimized for the signal.

3.1 R_p-violating Searches

Since the Higgs boson can decay into 3 and 4 leptons with little missing energy, searches for R-parity violation from the "LLE" operator may have a decent efficiency for selecting these

events. In particular, the CDF [20] and D0 [21] searches for this type of R-parity violation look for 3 charged leptons of the form lll or $ll\tau$ where $l = e$ or μ and there is no requirement on the flavors of l. Importantly, there is no cut on missing energy.

In the R-parity violating model the leptons are produced from the decay of the lightest neutralino through the gaugino-lepton-slepton and R-parity violating "LLE" operators. To set limits, two minimal supergravity models were considered where the parameters were chosen such that chargino-pair and chargino-neutralino production, rather than cascade decays of squarks and gluinos to neutralinos and charginos, were the main source of leptons. For operators involving only the first and second generation the limits on the production cross-section to these final states are quite strong, of the order 0.05 pb for D0, $\simeq 0.07 - 0.1$ pb for CDF, and are approximately independent of the mass of the lightest neutralino.

In direct production of the Higgs boson, charged leptons are produced in the cascade decay $h \to N_I N_I \to l + \cdots$. Since the decay of a single right-handed neutrino can produce at most two charged leptons, four charged leptons only occur if both right-handed neutrinos decay this way. Three charged leptons are produced if one right-handed neutrino decays to a single charged lepton and the other decays to two. From Table 2, the branching fraction for a right-handed neutrino to decay inclusively into jets and a charged lepton, summed over all lepton flavors including τ, is approximately 0.5. The branching fraction for it to decay into two charged leptons and missing energy, summed over all flavors, is approximately 0.2. The branching fraction for the two right-handed neutrinos to decay into three or more charged leptons, summed over all flavors, is then $2(1/2)(0.2) + (0.2)^2 = 0.24$.

Another source of charged leptons is if the Higgs boson is produced in association with a W or Z boson that decays leptonically. If $W \to l\nu$ occurs, then the two right-handed neutrino are required to produce at least two charged leptons, rather than three. If $Z \to l\bar{l}$, then the two right-handed neutrinos need only produce at least one charged lepton. One finds the branching fraction for two right-handed neutrinos to decay into at least two (one) charged leptons, summed over lepton flavors, is approximately 0.6(0.9).

In sum the rate for producing three or more charged leptons is

$$\begin{aligned}\sum_{\text{lepton flavor}} \sigma(p + \overline{p} \to h + Y \to l'l''l''' + X) &\simeq [2.5(0.24)\text{pb} + (0.3\text{pb})(0.2)(0.6) \\ &\quad +(0.2\text{pb})(0.06)(0.9)] \sum_I Br(h \to N_I N_I) \\ &\leq 0.63\text{pb} \end{aligned} \tag{3.1}$$

This result includes decays to τ leptons, which were needed to obtain a model-independent result. Since $l\tau\tau$ and $\tau\tau\tau$ final states were not included in either searches, this result overestimates the number of events eventually accepted by the searches. It should be interpreted as an upper bound that is independent of the neutrino mass parameters. I have also included in the first term events from associated Higgs production where the gauge bosons do not decay to charged leptons. Higgs production cross-sections are estimated as 2 pb for gluon-gluon fusion, 0.3 pb for associated W production, and 0.2 pb for associated Z production [22]. The

reader will notice that the estimate above is about an order of magnitude larger than the limits obtained for the R$-$parity violating models.

I will use these values to set limits on decays of the Higgs boson into right-handed neutrinos. To do this I will assume that the acceptances for the Higgs boson and R-parity violating decays are similar. In actuality, there are two factors which reduce the acceptance. Most of the signal events are from gluon-gluon fusion Higgs production. But since the Higgs boson decay products are soft (6 body decay), the acceptance for these events will be lower than in the 3-body neutralino decay into leptons (in the D0 search, the lightest neutralino mass is varied over $\approx 90 - 140$ GeV, which is comparable to the mass range for the Higgs boson considered here). A weaker limit is then expected, but a detailed study is needed. Events from associated production of the Higgs boson are not expected to have significantly different acceptance, since there is one charged lepton (or 2 if from Z decay) with high p_T. However, here the net cross section is smaller, about 0.03 pb which is below the limit set on the $R-$parity violating models.

The other condition reducing the acceptance is the following. Both experiments require that a number of the leptons be prompt. This means that the minimum separation between the lepton and the primary vertex projected onto the transverse plane must not exceed some value d_0. Specifically, for this analysis CDF requires that $d_0 < 0.2$ cm and $D0$ requires that $d_0 < 2$ cm. A right-handed neutrino with average decay length $D_I \gg O(d_0)$ will have a much lower acceptance, since most of the decays will have a minimum transverse distance exceeding d_0. However, since the decay occurs at random, the location of the secondary vertex forms a distribution and a small fraction of events will decay with a transverse distance $d_T < d_0$. To estimate that, I will crudely approximate the allowed region as a sphere of radius $R < O(d_0)$. This is a crude approximation, since actual detectors are tubular rather than spherical, and moreover the requirement is only on the transverse distance, not the physical distance. Then the probability that a right-handed neutrino with average decay length D_I decays within a sphere of size $R = d_0$ is approximately

$$P(d < d_0; D_I) = 1 - e^{-d_0/D_I} \ . \tag{3.2}$$

The average decay length D_I depends on the right-handed neutrino flavor and on its velocity β_I and boost γ_I. It is given by

$$D_I = \gamma_I \beta_I c \tau_I \ . \tag{3.3}$$

The estimate (3.1) for the inclusive production cross-section for three or more charged leptons is then modified to

$$\sum_{\text{lepton flavor}} \sigma(p + \overline{p} \to h + Y \to l' l'' l''' + X) \simeq \sum_I Br(h \to N_I N_I) \left[0.6 P(d < d_0; D_I)^2 \right. \\ \left. + 0.04 P(d < d_0; D_I) \right] \text{pb} \tag{3.4}$$

This result is maximized when all of the right-handed neutrinos have the same average decay length $D_I = D_*$ (see the discussion below (3.8)). Then

$$\sum_{\text{lepton flavor}} \sigma(p+\overline{p} \to h+Y \to l'l''l''' + X) \leq \big[0.6P(d<d_0;D_*)^2 + 0.04P(d<d_0;D_*)\big]\,\text{pb} \times \sum_I Br(h \to N_I N_I) \qquad (3.5)$$

The stronger limit is from D0, since they have a weaker cut on d_0. Assuming the signal here has the same acceptance as the signal in the R$-$parity violating models, then with $D_* \gtrsim 10$ cm and $\sum_I Br(h \to N_I N_I) = 1$, the production cross section is ≤ 0.03 pb which is less than but close to the $D0$ bound. The actual limit may be weaker than this estimate for several reasons. The acceptance for events from primary Higgs production may be lower, since the leptons produced in the Higgs boson decay are softer than those produced in the decay of the lightest neutralino. Events from associated Higgs production may not have a lower acceptance, but their rate, given by the second term in (3.5), is much smaller. And as mentioned above, this result overestimates the signal since it includes all the decays of right-handed neutrinos to τs.

In a specific model the production cross section may not saturate the model-independent upper bound derived above. For instance, if the average decay lengths are hierarchical then only one right-handed neutrino may have a significant contribution to (3.5). If so, the total rate is suppressed by the branching for the Higgs boson to decay into those right-handed neutrinos.

3.2 Inclusive Di-lepton Searches

CDF has an inclusive search for di-lepton events where the two leptons have the same charge [24]. Since the decay of right-handed neutrinos can produce leptons of the same sign, in principle the branching fraction for decays of the Higgs boson into right-handed neutrinos could be constrained by this search. These events occur from a number of processes and a model-independent upper bound on their production rate is derived below.

First consider primary production of the Higgs boson which then decays to right-handed neutrinos. A single right-handed neutrino can produce a single charged lepton and quarks with a high branching fraction. With a smaller branching fraction it can decay to two charged leptons and missing energy. Here though the two charged leptons have the opposite sign. Events with two charged leptons of the same-sign therefore requires that both right-handed neutrinos decay leptonically. From Table 2, the branching fraction for a right-handed neutrino to decay inclusively into jets and a charged lepton, summed over all charged lepton flavors (including τ), is approximately 0.5. The probability for the two right-handed neutrinos to decay into jets and two charged leptons of the same sign and summed over all 3 lepton flavors is therefore $(0.5)^2/2 = 1/8$. The branching fraction for a single right-handed neutrino to decay inclusively into two charged leptons, summed over charged lepton flavors, is approximately

0.21, so the branching fraction for two right-handed neutrinos to give three charged leptons, summed over all flavors, is $(0.5)(2)(0.21) = 0.21$. In this case two of the leptons will have the same charge. Finally, the branching fraction to obtain four charged leptons and missing energy, summed over all charged lepton flavors, is approximately $(0.2)^2 = 0.04$. In total, the branching fraction for two right-handed neutrinos to decay into at least two charged leptons of the same sign and summed over all lepton flavors is approximately 0.37. The rate for producing two *prompt* same-signed di-leptons from primary Higgs production and summed over all lepton flavors is then

$$\begin{aligned}\sum_{\text{lepton flavor}} \sigma(p+\overline{p} \to h+Y \to l^{\pm}l'^{\pm}+X)_{prmy} &= (2.5\text{pb})(0.37) \\ &\quad \times \sum_I Br(h \to N_I N_I)(P(d<d_0;D_I))^2 \\ &= 0.93\text{pb} \\ &\quad \times \sum_I Br(h \to N_I N_I)(P(d<d_0;D_I))^2\end{aligned} \tag{3.6}$$

The two factors of $P(d < d_0, D_I)$ estimates the probability that both of the right-handed neutrinos, each of average decay length D_I, decay inside the prompt signal region $d_0 <$ 2cm. By using a Higgs boson production cross-section of ≈ 2.5 pb for $m_h \simeq 115$ GeV I am also including associated Higgs boson production with an electroweak gauge boson (which contributes in total ≈ 0.5 pb to the Higgs boson production cross-section).

Events with same-signed leptons also occur with associated Higgs boson production, where one of the gauge bosons decays leptonically. The efficiency for events from associated production will be higher, since they have an energetic charged lepton, and if the lepton is from a W, significant missing energy. If $W \to l^{\pm}\nu(\overline{\nu})$ occurs, then only one right-handed neutrino is required to decay into a charged lepton of the same sign. If $Z \to l^- l^+$, the right-handed neutrino can decay to a charged lepton of either sign, although the rate for these events is small. The branching fractions for a right-handed neutrino to decay into a single charged lepton and summed over all lepton flavors is approximately 0.5, and to decay into two charged leptons and summed over all lepton flavors is 0.21. Next, the probability that at least one of the two right-handed neutrinos has a prompt decay is $2P(d < d_0; D_I)(1 - P(d < d_0; D_I))$.(Events where both right-handed neutrinos decay promptly was included in the previous result (3.6). To include it here would be to double count.) Then the contribution from associated production, where the electoweak gauge boson decays leptonically and only one right-handed neutrino decays promptly, is given by

$$\begin{aligned}\sum_{\text{lepton flavor}} \sigma(p+\overline{p} \to h+W/Z \to l^{\pm}l'^{\pm}+X) &= \Big[\frac{1}{2}(0.2)(0.3\text{pb})(0.5)+(0.2)(0.3\text{pb})(0.21)+ \\ &\quad +(0.06)(0.2\text{pb})(0.72)\Big]\end{aligned}$$

$$
\begin{aligned}
&\times \sum_I Br(h \to N_I N_I) \\
&\times 2(P(d < d_0; D_I))(1 - P(d < d_0; D_I)) \\
= {}& 0.072\text{pb} \\
&\times \sum_I Br(h \to N_I N_I) P(d < d_0; D_I) \\
&\quad \times (1 - P(d < d_0; D_I))
\end{aligned}
\tag{3.7}
$$

Then the total inclusive rate for producing two leptons of the same sign, summed over all three lepton flavors, is

$$
\begin{aligned}
\sum_{\text{lepton flavor}} \sigma(p + \overline{p} \to h \to l^{\pm} l'^{\pm} + X)_{tot} = {}& 0.86\text{pb} \sum_I Br(h \to N_I N_I)(P(d < d_0; D_I))^2 \\
&+ 0.072\text{pb} \sum_I Br(h \to N_I N_I) P(d < d_0; D_I)
\end{aligned}
\tag{3.8}
$$

In (3.8) decays to τ leptons were included to obtain a result independent of the Dirac matrix elements. It therefore overestimates the number of signal events accepted by the experimental analysis, since the search only looked for electrons and muons (including those from τ decays). Branching fractions for decays into charged leptons with hadronically decaying τ not included can be computed from Table 1, but are clearly model-dependent. The result of such a computation is bounded from above by (3.8).

Of the right-handed neutrinos, there will be one that has the shortest decay length and therefore largest probability factor, call it P_*. The other right-handed neutrinos will have $P(d < d_0; D_I) \leq P_*$. It is then clear that the above cross-section is maximized when all of the right-handed neutrinos have the same $P(d < d_0; D_I)$ or, that is, the same decay length. Then

$$
\sum_{\text{lepton flavor}} \sigma(h \to l^{\pm} l'^{\pm} + X)_{tot} \leq \left[0.86\text{pb}(P_*)^2 + 0.072\text{pb}P_*\right] \sum_I Br(h \to N_I N_I) \tag{3.9}
$$

I caution the reader that this is an upper bound that might not be saturated in specific models.

The CDF analysis [24] requires that the pseudorapidity of each lepton satisfy $|\eta| < 1.1$ and that the leading lepton has $E_T > 20$ GeV. These cuts will introduce an acceptance that is crudely 1/2 for each lepton, and 3/4 for the transverse energy cut, or ≈ 0.2 in total. From (3.9) with $D \approx 10$ cm and setting $\sum_I Br(h \to N_I N_I) = 1$, there are approximately 8 accepted events per 1 fb^{-1}. The analysis finds 44 events consistent with 33 expected background events. Decay lengths on the order of 10 cm do not appear to be excluded by this search, although a detailed studied is certainly warranted. A more careful analysis could probably

place interesting constraints for $D \leq O(5\text{cm}{-}10\text{cm})$ assuming $\sum_I Br(h \to N_I N_I) \simeq 1$, or $\sum_I Br(h \to N_I N_I) \lesssim O(10^{-2})$ for $D_I \ll O(\text{cm})$.

In specific models the upper bound (3.9) may not be saturated. For instance, if the decay lengths are hierarchical, then only one right-handed neutrino flavor may have a significant contribution to the sum (3.8). If so, the overall rate is suppressed by the branching ratio for the Higgs to decay into that state. In a minimal flavor violation scenario where the couplings (2.1) are unsuppressed, that branching ratio is approximately 1/3. Then with one decay length of approximately 10 cm there would only be roughly 3 accepted events per fb. If all of the decay lengths are larger than 10 cm then the acceptance decreases since there are fewer prompt events.

3.3 Long-lived Neutral Particles: Displaced Muons and Delayed Photons

The best strategy for finding the right-handed neutrinos is to look for displaced vertices. Standard Model backgrounds are practically nonexistent, so such events should be easy to find and interpret. Both CDF and D0 [18] have searches for a long-lived neutral particle decaying into specific channels, such as into a muon pair and neutrino [18], or into a photon with jets and missing energy [23], or to a physical Z boson [28]. Although these searches are optimized to constrain models of supersymmetry with R-parity violation or with low-energy gauge mediated supersymmetry breaking, their analysis should apply to the right-handed neutrino decays discussed here. While it shall be seen below that no limits yet exist, they are close. A reanalysis of their data after optimizing the search for the specific decay modes of the right-handed neutrinos could provide interesting limits.

Although there are several older searches, the best limits are from the D0 search [18] and the CDF search [23].

Limits on the model can be improved by optimizing the search to other decay modes having a higher branching ratio. For example, the right-handed neutrino has a branching ratio to a muon and quark pair that is about an order of magnitude larger than to a muon pair and neutrino. By searching for jets and a single lepton that reconstruct to a displaced vertex, the sensitivity to this model should in principle be improved.

3.3.1 Displaced Muon pairs

The DO search [18] looks for a neutral particle that travels a macroscopic distance and then decays into a neutrino and a pair of muons having opposite signs. In particular, they specifically look for a pair of opposite-signed muons that reconstruct to a secondary vertex. To have a high efficiency at reconstructing the muons, they restrict the analysis to decays that occur within the inner region of their central detector, or 5cm$< d <$ 20 cm. Right-handed neutrinos having decay lengths outside of this region will have a reduced acceptance and the limits will be weaker. The following discussion implicitly assumes that all three right-handed neutrinos have decay lengths in this region, although the formulae allow for the more general situation.

Each event will have two right-handed neutrinos pair-produced in the decay $h \rightarrow N_I N_I$. Moreover, at least one of these right-handed neutrinos must decay within the signal region into a muon pair, since both muons must come from the same vertex. The probability for that to happen is denoted by $P^{(I)}_{displ}$ and estimated below. The effective cross-section for producing di-muon pairs in the signal region is then

$$\sigma(\text{displaced } \mu^+\mu^- + X) = \sigma(p + \overline{p} \rightarrow h + X) \sum_I \left[Br(h \rightarrow N_I N_I) P^{(I)}_{displ} \right] . \tag{3.10}$$

Next I will estimate $P^{(I)}_{displ}$ by splitting it into two parts.

The probability of a right-handed neutrino decaying into $\mu^+\mu^-$ is flavor-dependent and denoted by $P^{(I)}_{\mu^+\mu^-}$. From Table 1 and using $|D_{JI}|^2 = |[m_D]_{IJ}|^2 M_I^{-1}$ this probability is given by

$$P^{(I)}_{\mu^+\mu^-} = \frac{|[m_D]_{I2}|^2}{[m_D m_D^\dagger]_{II}} \left(\frac{0.59}{N_{tot}} \right) + \left(\frac{0.13 c_Z}{N_{tot}} \right) \sum_{J \neq 2} \frac{|[m_D]_{IJ}|^2}{[m_D m_D^\dagger]_{II}} \tag{3.11}$$

The first term in the first line is due to interfering charged current and neutral currents, and the second set of terms is due to neutral currents. Using $N_{tot} \simeq 12$ and $c_Z \simeq 1$, an approximation to (3.11) is

$$P^{(I)}_{\mu^+\mu^-} \simeq 0.05 \frac{|[m_D]_{I2}|^2}{[m_D m_D^\dagger]_{II}} + 0.01 \sum_{J \neq 2} \frac{|[m_D]_{IJ}|^2}{[m_D m_D^\dagger]_{II}} \tag{3.12}$$

which is bounded from above by

$$P^{(I)}_{\mu^+\mu^-} \leq 0.05 \frac{1}{[m_D m_D^\dagger]_{II}} \left[|[m_D]_{I2}|^2 + \sum_{J \neq 2} |[m_D]_{IJ}|^2 \right] = 0.05 . \tag{3.13}$$

A lower bound is

$$P^{(I)}_{\mu^+\mu^-} \geq 0.01 \tag{3.14}$$

and is saturated only if $[m_D]_{I2} = 0$. Saturation of the upper bound occurs only if $[m_D]_{IJ} = 0$ for $J \neq 2$. This condition is not possible to physically realize since it requires the Dirac mass matrix to be of rank 1 and therefore implies that two active neutrinos are massless. Another way to say that is that this condition requires that only ν_2 has mass mixing with the right-handed neutrinos, and therefore ν_1 and ν_3 are decoupled and remain massless. Nonetheless, it is useful to have a model-independent upper bound, albeit not possible to saturate given the data on neutrino masses.

Next, I approximate the combined probability that at least one right-handed neutrino in the pair decays in the signal region into a muon pair as

$$\begin{aligned} \mathcal{P}^{(I)}_{displ} &= 2P_I(1 - P_I) P^{(I)}_{\mu^+\mu^-} + 2P_I^2 P^{(I)}_{\mu^+\mu^-}(1 - P^{(I)}_{\mu^+\mu^-}) + P_I^2 (P^{(I)}_{\mu^+\mu^-})^2 \\ &= 2P_I(1 - P_I) P^{(I)}_{\mu^+\mu^-} + P_I^2 P^{(I)}_{\mu^+\mu^-}(2 - P^{(I)}_{\mu^+\mu^-}) . \end{aligned} \tag{3.15}$$

Here P_I is the probability that a right-handed neutrino N_I with average decay length D_I decays within the signal region. As in previous sections, to estimate that I treat the signal region as a sphere - rather than a tube - extending from $R = 5$ cm out to $R = 20$ cm and approximate $P_I \equiv P(5\text{cm}, 20\text{ cm}; D_I)$ where

$$P(R_1, R_2; D) \equiv e^{-R_1/D} - e^{-R_2/D} \tag{3.16}$$

is the probability that a right-handed neutrino of average decay length D decays a distance between $r = R_1$ and $r = R_2$ from the primary vertex. The first term in $P^{(I)}_{displ}$ is the probability that one right-handed neutrino decays in the signal region into a muon pair, the second term is the probability that both decay in the signal region, but only one decays into a muon pair, and the third term describes the case where both decay in the signal region into muon pairs.

An upper bound on (3.10) can be obtained by first noting from (3.13) that $P^{(I)}_{\mu^+\mu^-} < 0.05$. Then $P^{(I)}_{displ} \leq 0.1P_I(1-P_I) + 0.1P_I^2 < 0.07$ with the maximum occurring around $D_I \simeq 10$cm. Then with $\sigma(p + \overline{p} \rightarrow h + X) \leq 2.5$ pb, $\sum_I BR(h \rightarrow N_I N_I) \leq 1$ and setting $D_I \simeq$10 cm for all three right-handed neutrinos to maximize the rate, one obtains the model-independent bound

$$\sigma(\text{displaced } \mu^+\mu^- + X) < 0.175 \text{ pb}\ . \tag{3.17}$$

As discussed below (3.13), it is not possible to saturate this bound with any model consistent with what is known about the active neutrino masses. However, the lower bound 3.14 suggests that values might be only a factor of a few smaller in models satisfying the assumptions listed above.

For comparison, the limits from D0 [18] vary considerably, with their best limit given by 0.14 pb which is slightly below the upper bound (3.17). That the model-independent bound is slightly larger than the best experimental limit probably does not restrict any of the parameter space, since as I have already stressed, the model-independent bound cannot be saturated with any realistic model since it would require two active neutrinos to be massless. Yet that these upper bounds are so close to one another warrants a more careful analysis than described here, in order to more properly determine the acceptance of the signal in the signal region. A meaningful lower bound to the signal rate cannot be obtained since the rate decreases rapidly once the average decay lengths lie outside of the signal region.

3.3.2 Displaced Photons

The CDF search [23] for displaced photons looks at inclusive events with a photon having $E^T > 25$ GeV, at least one jet with $E_T > 30$ GeV, and missing energy greater than approximately 30 GeV [23]. In principle long-lived decays of right-handed neutrinos could be constrained by this analysis, since it has a rare decay $N_I \rightarrow \gamma$+missing energy. The required jet can be produced from either the decay of the other right-handed neutrino in the event, or from a hadronically decaying electroweak gauge boson produced in association with the Higgs boson.

Photons were required to be "out-of-time" with the primary events by approximately 2-10ns [23]. Whether this occurs in a model depends on the right-handed neutrino's average decay length and whether it is non-relativistic. Both conditions are required [23]: a non-relativistic right-handed neutrino that decays promptly ($\tau \ll O(\mathrm{ns})$) will not produce a delayed photon; and a relativistic right-handed neutrino with a long average decay length will produce on average a photon moving in the same direction, which arrives at the calorimeter simultaneously with other relativistic particles in the event. Values of the underlying parameters satisfying these conditions certainly occur. For a lifetime of $O(2\mathrm{ns} - 10\mathrm{ns})$ corresponds to a $c\tau$ of approximately $O(60cm - 3m)$, which occurs for a wide range of neutrino mass parameters (see Eqn. 3). Non-relativistic right-handed neutrinos, with say $\beta \lesssim 0.6$, are more difficult to obtain, but they do occur for masses $M_I \gtrsim 0.4 m_h$. The actual region of parameters having a decent acceptance will not be identified, but it is clear that it is not tiny. The comments below are intended to apply to this region.

If there are other energetic, promptly decaying particles occurring in the event, then they will arrive at the detector before the photon. For these class of events the distribution of the time delay for the photon is highly asymmetric, as it does not extend into regions with negative values of the time delay Δt. This feature was utilized in the CDF search to measure their background to the signal. There the background to the signal, occurring from beam-related events and Standard Model processes, was obtained by extrapolating a measurement of the background in the region $\Delta t < 0$, where no signal is expected, into the region $\Delta t > 0$ were a signal is expected.

This feature of the analysis raises an issue for signal events occurring from primary Higgs production. For here all the 'hard events' are delayed. On average the decay of the two right-handed neutrinos will be out-of-time by $O(1/\Gamma)$ with respect to each other. But the jet used in the event selection is produced from the decay of a right-handed neutrino, which is itself delayed. In half of the events then, the jet will be delayed with respect to the photon. The distribution of the photon's time delay with respect to the arrival time of the jet will then be symmetric about zero. How well the CDF analysis can constrain such signal events is important to determine, since the search is beginning to cut into an interesting region of parameter space (see below). Certainly a limit can be set, since only a few background events were observed in the $\Delta t < 0$ region (from Fig. 1 of [23]).

This issue does not arise in associated production of the Higgs boson, since the photon from the decay of one of the right-handed neutrinos is delayed with respect to the particles produced from the prompt decay of the electroweak gauge boson.

The effective production cross-section for the signal events is

$$\begin{aligned}\sigma(\text{displaced}\gamma + \text{jets+missing energy}) = {} & \sigma(p + \overline{p} \to h + X) \sum_I Br(h \to N_I N_I) \\ & \times 2 Br(N_I \to \gamma + \text{missing energy}) \\ & \times Br(N_I \to \text{jets} + X) \end{aligned} \tag{3.18}$$

As discussed in Section 2.2.2, the branching fraction for the radiative decay is approximately

$7 \times 10^{-5}|c_M|^2$ and is independent of the Dirac matrix elements if minimal flavor violation is assumed. c_M is a linear combination of a calculable one-loop contribution and an unknown coefficient $c_M^{(6)}$ which is determined by the coefficients of the two dimension 6 operators that contribute to this decay. While this nominal value for the branching fraction is too tiny to be relevant here, it increases to $0.007|c_M^{(6)}|^2(\mathrm{TeV}/\Lambda)^4$ for $\Lambda \lesssim 3$ TeV.

Missing energy in the event occurs from the radiative decay producing on average approximately $m_h/4$. Additional missing energy can be obtained from the hadronic decay of the second right-handed neutrino in the event, or from a leptonic decay of a W, if present. From Table 2 the branching fractions for a right-handed neutrino to decay into light quarks or heavy quarks and missing energy totals to 0.21 and is model-independent. The branching fraction for it to decay into jets and any charged lepton is approximately 0.5 and also model-independent. Although these latter decays do not produce any missing energy, I will be conservative and include them in my estimate below. I note that if $m_h \gtrsim 120$ GeV then such events can produce on average enough missing energy from the radiative decay alone to pass the missing energy cut. Then

$$\begin{aligned}\sigma(\text{displaced}\gamma+\text{jets+missing energy}) &= 2.5\text{pb} \times 2 \times 7 \times 10^{-3}|c_M^{(6)}|^2 \left(\frac{\mathrm{TeV}}{\Lambda}\right)^4 \times 0.71 \\ &\quad \times \sum_I Br(h \rightarrow N_I N_I) \\ &= 0.025\text{pb}|c_M^{(6)}|^2 \left(\frac{\mathrm{TeV}}{\Lambda}\right)^4 \sum_I Br(h \rightarrow N_I N_I) \quad (3.19)\end{aligned}$$

where both primary and associated Higgs production have been included. The "2" counts the two possibilities for a radiative decay from one of the two right-handed neutrinos. With 100% acceptance and 570pb^{-1} of data and $\sum_I Br(h \rightarrow N_I N_I) = 1$ there would have been 14 events produced for these fiducial values. The CDF analysis imposed a number of cuts that reduces the acceptance of the signal. For the gauge-mediated supersymmetry breaking models analyzed in [23] the acceptance is listed in their Table 1 as 23%. Assuming a similar acceptance here implies 3 accepted events for the fiducial values above. Since only 1 event above background was seen, a mild suppression of $|c_M^{(6)}| \simeq 0.6$ results in only 1 accepted event. With this value, approximately 5 events would have been produced. This estimate of the limit on $|c_M^{(6)}|$ is consistent with a conclusion of [23] that their analysis is general enough to exclude any model producing more than 5.5 events containing jets, missing energy and a delayed photon with time-delay of $O(2\text{ns} - 10\text{ns})$. Increasing Λ by a factor of 2 decreases the rate by a factor of 16, so no limit can be set for $|c_M^{(6)}| = 1$, $\Lambda \simeq 2$ TeV (≈ 1 event would have been produced). In interpreting these conclusions, it is important to recall that they only apply to those values of the mass parameters which result in non-relativistic right-handed neutrinos. Yet it is interesting that the search for delayed photons is beginning to exclude $\Lambda \simeq 1$ TeV if the dominant decay of the Higgs boson is to non-relativistic right-handed neutrinos.

Further analysis of these decays to optimize the experimental search is certainly warranted. Moreover, it is important to determine the sensitivity of the search strategy to

delayed photons occurring from primary Higgs production, for the reasons already discussed above. If this is indeed an issue, then the constraint obtained above may have been too cautious.

3.4 Inclusive Searches for Anomalous Photons

Signature-based searches [26, 27]involving photons are of several types. There are inclusive searches for $\gamma\gamma+X$ where $X=\gamma$, missing energy, or a prompt lepton. There are also inclusive searches for γl and γll. These searches may potentially constrain the model, since photons may be produced from both real and fake physics sources.

As discussed in Section 2.2.2, additional operators at dimension 6 introduce a small branching fraction for right-handed neutrinos to decay

$$N_I \rightarrow \gamma\nu_L, \gamma\overline{\nu}_L \ . \tag{3.20}$$

Since the right-handed neutrinos are displaced a macroscopic distance on average, these photons will not point back to the primary vertex. From (2.21), the inclusive branching fraction for this decay is approximately

$$\begin{aligned} Br(N_I \rightarrow \gamma + \text{missing energy}) &= 7\times 10^{-5}|c_M|^2 \\ &\simeq |c_M^{(6)}|^2 \left(\frac{\text{TeV}}{\Lambda}\right)^4 \ . \end{aligned} \tag{3.21}$$

The last relation occurs for $\Lambda \lesssim 4\pi v \simeq 3$ TeV, which will be assumed for the remainder of this subsection.

To estimate the event rate, note that summing over the three dominant production mechanisms for the Higgs boson gives $\sigma_h \simeq 2.5$ pb for $m_h \simeq 115$ GeV . Tnen the inclusive rate to produce a photon from one of the two right-handed neutrinos is

$$\sigma(h \rightarrow \gamma + \text{missing energy} + X) \simeq 35 \sum_I Br(h \rightarrow N_I N_I)|c_M^{(6)}|^2 (\text{TeV}/\Lambda)^4 \text{ fb} \tag{3.22}$$

where X is not yet specified. The energy of the photon is $\approx m_h/4 \approx 25$ GeV and the missing energy is $\approx m_h/4 \approx 25$ GeV plus whatever is lost in X.

The largest rate occurs from primary Higgs production. From Table 2 the second right-handed neutrino will decay into charged leptons plus jets 50% of the time, and into two charged leptons and missing energy about 24% of the time. As these leptons must be prompt in order to be included in the analysis, such events are suppressed by a factor of $P(d<d_0=2\text{ cm}; D_I)$. Leptons from decays of a W produced in association with the Higgs are prompt, but the overall rate is lower. Adding both production mechanisms gives

$$\begin{aligned} \sigma(h \rightarrow \gamma + l + \text{missing energy} + X) \simeq & \sum_I Br(h \rightarrow N_I N_I) \left[35 P(d<d_0; D_I)(0.75)\right. \\ & \left. +(0.007)(300)(0.2)\right] \\ & \times |c_M^{(6)}|^2 (\text{TeV}/\Lambda)^4 \text{ fb} \end{aligned}$$

$$= \sum_I Br(h \to N_I N_I)\,[26P(d < d_0; D_I) + 0.4]$$
$$\times |c_M^{(6)}|^2 (\mathrm{TeV}/\Lambda)^4 \text{ fb} \tag{3.23}$$

With $D_I > 10$ cm, $P(d < 2 \text{ cm}; D_I) \lesssim 0.2$ and setting $\sum_I Br(h \to N_I N_I) = 1$ gives approximately $6|c_M^{(6)}|^2 (\mathrm{TeV}/\Lambda)^4$ events produced per fb^{-1}. The number of events detected is reduced by the acceptance which I have not included here. From inspecting Figures 2,3 and 4 from the CDF Public Note [27] giving distributions of lepton plus photon plus missing energy events, a handful (≈ 13 when combined with (3.29) below) of such events do not appear to be excluded, although a more careful analysis is warranted. Limits can probably be set for smaller values of the average decay length. In particular, for $D_I \ll O(5\text{cm})$ limits $\sum_I Br(h \to N_I N_I) \lesssim 0.1$ could probably be set assuming $c_M^{(6)} = 1$ and $\Lambda = 1$TeV.

In principle one might consider signal events containing two photons and missing energy $\simeq m_h/2 \simeq 60$ GeV. Then

$$\sigma(h \to \gamma\gamma + \text{missing energy} + X) \simeq 0.2 \sum_I Br(h \to N_I N_I) |c_M^{(6)}|^2 (\mathrm{TeV}/\Lambda)^4 \text{ fb} \tag{3.24}$$

which is unfortunately too small.

A potential source for fake photons occurs if a right-handed neutrino decays inside the electromagnetic calorimeter into a final state containing an electron. These electrons will be contained in the calorimeters and so will not produce a track. How this electron is identified depends on the analysis. Several Tevatron analyses explicitly state that a fully contained energy deposition in the electromagnetic calorimeter without an associated track will be identified as a photon. While such an electron is in principle distinguishable from a photon and might be misidentified as a neutral hadron, to be conservative in ascertaining whether these searches may have "caught" some of the signal events, I will assume that some of these electrons are misidentified as photons. This fraction will be denoted by $f_{\gamma|e}$ and I will not try to estimate it.

The electromagnetic calorimeter is shaped like a cylinder and not at a fixed distance from the primary vertex. It is therefore sensitive to a range of average decay lengths. Since the calorimeter is not big, the fraction of decays occurring in that region is suppressed. The CDF calorimeter has a width of about 30cm and extends from approximately $R_1 = 1.7$m out to approximately $R_2 = 2$m. The fraction of events decaying in this volume depends sensitively on the geometry of the calorimeter and requires a detector simulation. To obtain a rough estimate, I will crudely approximate the electromagnetic calorimeter as a hollow sphere with radius extending between 1.7m and 2m. Then the fraction of right-handed neutrinos of average decay length D_I that decay in this region is

$$P_\gamma(D_I) \equiv P(R_1 < d < R_2; D_I) \; . \tag{3.25}$$

Let me begin with the "$\gamma\gamma$"+missing energy events from primary Higgs production. Events with missing energy and two fake photons occur from $h \to N_I N_I$ where either one or

both right-handed neutrinos decay inside the electromagnetic calorimeter. If only one decays inside, then $N_I \to e^+e^-$+missing energy with branching fraction $P^{(I)}_{e^+e^-}$. One can show $P^{(I)}_{e^+e^-} \leq 0.05$ using reasoning similar to that found in Section 3.3.1 to derive an identical upper bound on decays to muon pairs. For decays to electrons, saturation of the bound requires $[m_D]_{IJ} = 0$ for $J \neq 1$ which as discussed in Section 3.3.1 is unrealistic. If both right-handed neutrinos decay in the electromagnetic calorimeter then the majority of the fake events occur when each decays to an electron. Although the dominant decay of N_I is into charged leptons and light quarks and occurs with branching fraction of 1/2, both right-handed neutrinos cannot decay into this channel since there is no missing energy. One N_I must decay to $e + l$+missing energy, whose branching ratio $P^{(I)}_{el}$ is less than 0.24. Then

$$\begin{aligned}\sigma(h \to \text{``}\gamma\gamma\text{''} + X) &= 2.5\text{pb} \sum_I Br(h \to N_I N_I) \\ &\quad \times f^2_{\gamma|e} \left[2P_\gamma(D_I)(1 - P_\gamma(D_I))P^{(I)}_{e^+e^-} + 2P^{(I)}_{eqq} P^{(I)}_{el} (P_\gamma(D_I))^2\right] \quad (3.26)\end{aligned}$$

Varying D_I to maximize the signal gives $P_\gamma(D_I) = P(1.7 < d < 2; D_I) \leq 0.06$ and $2P_\gamma(D_I)(1-P_\gamma(D_I)) \leq 0.12$. Using $P^{(I)}_{eqq} \leq 0.5$ and $P_{el} \leq 0.24$ one obtains finds

$$\begin{aligned}\sigma(h \to \text{``}\gamma\gamma\text{''} + X) &< 2.5\text{pb} f^2_{\gamma|e} \sum_I Br(h \to N_I N_I) \left[(0.12)(0.05) + 0.25(0.06)^2\right] \\ &= 0.017\text{pb} f^2_{\gamma|e} \sum_I Br(h \to N_I N_I) \quad (3.27)\end{aligned}$$

A more accurate estimate of the rate can be determined in specific models, but it will depend on the Dirac mass matrix elements. It will be bounded by the result above, but will probably only be a factor of a few smaller.

The number of events detected is reduced from the above result by several factors. First, there is the unknown probability $f_{\gamma|e}$ of misidentifying a trackless electron produced in the electromagnetic calorimeter as a photon. Next, the CDF analysis [26] requires that both photons have pseduo-rapidity $|\eta| < 1$, which lowers the acceptance. Finally, the acceptance is further reduced because the events described above have little missing energy, since most only have one neutrino that on averages has $E \simeq m_h/6$. Assuming 100% acceptance for the misidentification of electrons ($f_{\gamma|e} = 1$) and $\sum_I Br(h \to N_I N_I) = 1$, there are less than $\approx 17\text{fb}^{-1}$ produced events. Inspecting their Figures 1 and 2, in their control region they expected 115 events with missing $E_T > 20$ GeV, and observed 126. In their signal region they expected 174 and saw 196 events. So no constraint currently exists from this channel.

Events with a larger amount of missing energy require associated production of hW or hZ, where the gauge bosons decay to neutrinos. Then the analysis above mostly carries over, except that now if both right-handed neutrinos decay inside the electromagnetic calorimeter both can decay to a charged lepton and quarks, since no missing energy is needed from the right-handed neutrinos. Then with $Br(N_I \to e + X) \equiv P^{(I)}_e \leq 0.7$,

$$\sigma(h \to \text{``}\gamma\gamma\text{''} + X)_{asscprod} = (0.3)(0.3) f^2_{\gamma|e}\text{pb} \sum_I Br(h \to N_I N_I)$$

$$\times \left[2P_\gamma(D_I)(1-P_\gamma(D_I))P^{(I)}_{e^+e^-} + P^{(I)}_e P^{(I)}_e (P_\gamma(D_I))^2\right]$$
$$\leq 0.7\text{fb} f^2_{\gamma|e} \sum_I Br(h \to N_I N_I) \tag{3.28}$$

Next consider $\gamma l{+}X$ events where an electron produced in the electromagnetic calorimeter is faking a photon. The lepton must be prompt and can be produced from the other right-handed neutrino decaying leptonically in the acceptance region $d < d_0 \simeq 2$ cm, or from the leptonic decay of a W. The former events are suppressed by the additional factor of $P(d < 2\text{ cm}; D_I)$, and the latter events by the lower production cross-section. In total

$$\begin{aligned}\sigma(p+\overline{p} \to h+Y \to l\text{``}\gamma\text{''}+X) &= \sum_I Br(h \to N_I N_I)\left[(2)2.5\text{pb}P(d<d_0, D_I)P_\gamma(D_I)P^{(I)}_e f_\gamma\right.\\ &\quad \left. +0.3\text{pb}(0.2)P_\gamma(D_I)P^{(I)}_e f_\gamma\right]\\ &< 7 f_\gamma \text{ fb} \sum_I Br(h \to N_I N_I)\end{aligned}$$

where $P(d < 2\text{cm}, D_I) \times P_\gamma(D_I) < 1.2 \times 10^{-3}$, $P_\gamma(D_I) < 0.06$ and $P^{(I)}_e < 0.75$ have been used. With 1 fb^{-1} this leads to a handful of events produced. Inspecting Figures 2 and 3 in [27] giving the distribution for events in the $e\gamma$ and $\mu\gamma$ missing energy samples, it appears that the addition of a handful of events produced (when combined with (3.23)) in each channel is not excluded.

4. Beyond Minimal Flavor Violation

The minimal flavor violation hypothesis [13] has been used throughout since dimension 6 operators involving *quarks* or *leptons* must be suppressed in order to be consistent with flavor changing neutral current processes. This hypothesis seems too restrictive however, for operators that involve only right-handed neutrinos, since they are more poorly constrained. In particular, there are no experimental constraints on the operator

$$\frac{c_{IJ}}{\Lambda} N_I N_J H^\dagger H \tag{4.1}$$

(assuming $\Lambda > O(m_R)$ for consistency of the effective theory).

With minimal flavor violation, the couplings c_{IJ} are aligned with the right-handed neutrino masses m_R, so that they are simultaneously diagonalizable. In that case the couplings to the Higgs boson are diagonal in the right-handed neutrino mass basis.

In a more general context though with arbitrary c_{IJ} not aligned with the right-handed neutrino masses, non-universal couplings can occur and they have interesting consequences. For now the couplings to the Higgs boson are

$$\frac{v}{\Lambda} c'_{IJ} N_I N_J h \tag{4.2}$$

with $c'_{IJ} \neq \delta_{IJ}$. Now heavier right-handed neutrinos can decay into lighter right-handed neutrinos through an off-shell h. This leads to the following new decays [3]

$$N_I \to N_J h^* \to N_J b\bar{b} \tag{4.3}$$

where $M_I > M_J$. These decays are to a two-body final state and are only suppressed by the Higgs boson mass and the bottom Yukawa coupling. They will dominate over the charged current and neutral current decay processes for $c'\lambda_b \gtrsim O(\lambda_\nu)$ which occurs for a wide range of c''s.

One then has the following scenario. The heavier two right-handed neutrinos can decay to a bottom quark pair and a lighter right-handed neutrino. Depending on the relative sizes of the couplings c', the heaviest right-handed neutrino (N_H) may decay preferentially to the second heaviest right-handed neutrino (N_M), which then decays to the lightest right-handed neutrino (N_L) and another bottom quark pair. The lightest right-handed neutrino is kinematic forbidden to have these decays, so it will decay through the charged and neutral current processes. One could then have

$$\begin{aligned} N_H &\to 2N_L + b\bar{b} + b\bar{b} \\ &\to l + qq' + b\bar{b} + b\bar{b} \end{aligned} \tag{4.4}$$

which has 7 particles in the final state. The other decay process $N_H \to N_L + b\bar{b}$ could occur as well.

The couplings of the Higgs boson are now non-universal, so decays to all 6 possible pairing $N_I N_J$ occur. Decays of the Higgs boson into right-handed neutrino can produce anywhere from 4 and 6 particles in the case of Higgs boson decays to $N_L N_L$, to up to 14 particles in the case of decays to $N_H N_H$.

The phenomenology of this scenario is quite rich and the analysis of the experimental constraints somewhat more involved from what I considered in previous Sections. Because of the large number of b quarks in these events they probably would have been seen at LEP, so a limit of $m_h \gtrsim 114$ GeV can probably be set. The analyzes of the searches considered in previous Sections can be carried over to Higgs boson decays to $N_L N_L$, since the decays of the lightest right-handed neutrinos remains unchanged. However, the rates will be reduced by the branching fraction for the Higgs boson to decay exclusively to the two lightest right-handed neutrinos. The large number particles and of bottom quarks in the other decays will make it hard to find at hadron colliders.

5. Conclusion

I have discussed experimental limits on Higgs boson decays into right-handed neutrinos within the context of minimal flavor violation. These decays are interesting, since there is a wide

[3]The author thanks Scott Thomas for this observation.

range of scales Λ and Higgs mass in which they can be the dominant decay mode of the Higgs boson. To set precise limits however requires determining the acceptances for these decays, and exploring the parameter space of the model, and I have done neither. The parameter space for these models is large, having in addition 18 neutrino parameters. Setting limits is therefore model-dependent, since branching fractions for right-handed neutrinos decaying into specific final states are highly model-dependent. Model-independent branching fractions for inclusive decays can be derived (Table 2) and are quite useful in obtaining model-independent upper bounds to the number of events produced. I used these upper bounds to infer limits that might occur in an actual model, although that overestimates the rate.

Quite remarkably, much of that parameter spaces does not appear to be excluded. Dominant decays of the Higgs boson into right-handed neutrinos having average decay lengths larger than $O(10cm)$ appear to be allowed. Constraints from a number of searches do exist for average right-handed neutrino decays lengths of $O(2-5cm)$ and smaller, and probably require that the branching fraction for the Higgs boson to decay into such states is $O(10^{-2}-10^{-1})$. That the branching ratios are that small may naturally occur in specific models. For a shorter average decay length requires that the right-handed neutrino is heavier, with average decay lengths less than $O(cm)$ requiring that it be heavier than the W. Decays of the Higgs boson into these states may naturally be small because: i) the Higgs boson is light and this channel is not kinematically accessible, although decays to other right-handed neutrino pairs may be; or ii) the Higgs boson can decay into it, but then it may be heavy enough to decay into electroweak gauge bosons, in which case the branching fraction of the Higgs boson to decay into right-handed neutrinos is naturally small enough for $\Lambda \simeq O(5\text{TeV}-7\text{TeV})$ (compare Figure 3 with Figure 2a.). Current searches are beginning to probe regions of decay lengths larger than $O(10cm)$. These findings justify further studies to better determine the acceptances for these decays, and to optimize search strategies.

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A. Minimal Flavor Violation and Dimension 6 Operators

A predictive framework for the flavor structure of the higher dimension operators is provided by the minimal flavor violation hypothesis [13, 14, 15]. This hypothesis postulates a flavor symmetry assumed to be broken by a minimal set of non-dynamical fields or spurions, whose vevs determine the renormalizable Yukawa couplings and masses that violate the flavor symmetry. With the assumption of a single spurion breaking each flavor group, the flavor

structure of the higher dimension operators is fixed in terms of appropriate powers of the spurions, or in other words, in terms of the low-energy Yukawa couplings. Limits on operators in the quark sector are $5-10$ TeV [14], but weak in the lepton sector unless the neutrinos couplings are not much less than order unity [15][16].

If the assumption of minimality is relaxed then limits on the scale of these operators would be much higher. For if there are multiple spurions breaking the same flavor group, the linear combination of spurions appearing in a higher dimension operator does not have to be the same combination determining the Yukawa couplings. If so, the coefficients of the higher dimension operator cannot be expressed in terms of the Yukawa couplings. This is dangerous, since barring accidental cancelations, the misalignment between the higher dimension operators and the Yukawa couplings leads to large flavor violation. To avoid this potential problem, a minimal field content will be assumed [4].

The flavor symmetry in the lepton sector is taken to be

$$G_N \times SU(3)_L \times SU(3)_{e^c} \times U(1) \tag{A.1}$$

where $U(1)$ is the usual overall lepton number acting on the Standard Model leptons. With right-handed neutrinos present there is an ambiguity over what flavor group to choose for the right-handed neutrinos, and what charge to assign them under the $U(1)$. In fact, since there is always an overall lepton number symmetry unless both the Majorana masses and the neutrino coupling are non-vanishing, there is a maximum of two such $U(1)$ symmetries.

Two possibilities are considered for the flavor group of the right-handed neutrinos:

$$G_N = SU(3) \times U(1)' \text{ or } SO(3)\ . \tag{A.2}$$

The former choice corresponds to the maximal flavor group, whereas the latter is chosen to allow for a large coupling for the operator (2.1), shown below. The fields transform under the flavor group $SU(3) \times SU(3)_L \times SU(3)_{e^c} \times U(1)' \times U(1)$ as

$$N \rightarrow (\mathbf{3},\mathbf{1},\mathbf{1})_{(\mathbf{1},\mathbf{0})} \tag{A.3}$$
$$L \rightarrow (\mathbf{1},\mathbf{3},\mathbf{1})_{(\mathbf{-1},\mathbf{1})} \tag{A.4}$$
$$e^c \rightarrow (\mathbf{1},\mathbf{1},\mathbf{3})_{(\mathbf{1},\mathbf{-1})}\ . \tag{A.5}$$

Thus $U(1)'$ is a lepton number acting on the right-handed neutrinos and Standard Model leptons and is broken only by the Majorana masses. $U(1)$ is a lepton number acting only on the Standard Model leptons and is only broken by the neutrino couplings. Promoting the masses and Yukawa couplings of the theory to spurions gives for $G_N = SU(3) \times U(1)'$,

$$\lambda_\nu \rightarrow (\overline{\mathbf{3}},\overline{\mathbf{3}},\mathbf{1})_{(\mathbf{0},\mathbf{-1})} \tag{A.6}$$
$$\lambda_l \rightarrow (\mathbf{1},\overline{\mathbf{3}},\overline{\mathbf{3}})_{(\mathbf{0},\mathbf{0})} \tag{A.7}$$
$$m_R \rightarrow (\overline{\mathbf{6}},\mathbf{1},\mathbf{1})_{(\mathbf{-2},\mathbf{0})}\ . \tag{A.8}$$

[4] Multiple fields generating right-handed neutrinos masses are probably allowed, but not explored here.

For $G_N = SO(3)$ there are several differences. First, the $\overline{\mathbf{3}}$'s of $SU(3)$ simply become $\mathbf{3}$'s of $SO(3)$. Next, the $U(1)$ charge assignments remain but there is no $U(1)'$ symmetry. Finally, a minimal field content is assumed throughout, implying that for $G_N = SO(3)$ $m_R \sim \mathbf{6}$ is real.

As I discuss in [8], the sizes of the coefficents of the dimension 5 operators can now be estimated. There are several operators at dimension 5, but only (2.1) is potentially relevant to decays of the Higgs boson or right-handed neutrinos. All other operators are suppressed by powers of the neutrino couplings relative to the leading order decays. For (2.1) one finds that its coefficients depend on the choice of flavor group. This isn't surprising, since the operator $N_I N_I$ violates $SU(3)$, but preserves an $SO(3)$ symmetry. One finds

$$\begin{aligned} G_N = SU(3) \times U(1)' : c \sim a_1 \frac{m_R}{\Lambda} + a_2 \frac{m_R \mathrm{Tr}[m_R^\dagger m_R]}{\Lambda^2} + \cdots \\ G_N = SO(3) : c \sim \mathbf{1} + b_1 \frac{m_R}{\Lambda} + b_2 \frac{m_R \cdot m_R}{\Lambda^2} + \cdots + d_1 \lambda_\nu \lambda_\nu^\dagger + \cdots \end{aligned} \tag{A.9}$$

where $\cdots$ denotes higher powers in m_R and $\lambda_\nu \lambda_\nu^\dagger$. Comparing the expressions in (A.9), the only important difference between the two is that $\mathbf{1}$ is invariant under $SO(3)$, but not under $SU(3)$ or $U(1)'$. This is a key difference that has important consequences for the decay rate of the Higgs boson into right-handed neutrinos, as can be seen from Figures 1 and 2, and discussed further in [8].

At dimension 6 there are a number of operators. Here I focus on those involving at least one right-handed neutrino. One finds that to leading order in the Yukawa couplings

$$\mathcal{O}_1^{(6)} = g_1 c_1^{(6)} (N \lambda_\nu \sigma^{\mu\nu} L) \widetilde{H} B_{\mu\nu} \tag{A.10}$$
$$\mathcal{O}_2^{(6)} = g_2 c_2^{(6)} (N \lambda_\nu \sigma^{\mu\nu} \widetilde{H}^\dagger \tau^a L) W^a_{\mu\nu} \tag{A.11}$$
$$\mathcal{O}_3^{(6)} = c_3^{(6)} (N \lambda_\nu L)(L \lambda_l e^c) \tag{A.12}$$
$$\mathcal{O}_4^{(6)} = c_5^{(6)} (\overline{N} \lambda_\nu^* \lambda_l \gamma^\mu e^c)(H D_\mu H) \tag{A.13}$$
$$\mathcal{O}_5^{(6)} = c_6^{(6)} \partial^\mu N \lambda_\nu L D_\mu \tilde{H} \tag{A.14}$$
$$\mathcal{O}_6^{(6)} = c_7^{(6)} (N \lambda_\nu L \widetilde{H})(H^\dagger H) \tag{A.15}$$
$$\mathcal{O}_7^{(6)} = c_8^{(6)} \overline{N} \lambda_\nu^* \gamma^\mu \lambda_\nu N H^\dagger D_\mu H \tag{A.16}$$

where all of the $c_i^{(6)}$'s are numbers and not matrix-valued. Also, $\widetilde{H} \equiv i\tau_2 H^*$. Other operators are equivalent to those listed above either by an integration by parts in the action, or by using the free equations of motion.

Most of these operators introduce new decay modes for the right-handed neutrinos. But with the exception of a couple of operators discussed below, their amplitudes are parametrically suppressed by additional powers of Yukawa couplings compared to the previously discussed amplitudes created by mass mixing of the right-handed neutrinos with the left-handed neutrinos. For recall that the latter amplitudes are $O(\lambda_\nu)$. By inspection, all of the dimension 6 operators listed above contain one factor of the neutrino Yukawa coupling, a consequence of

the fact that the $U(1)$ lepton number is only violated by the neutrino couplings. Some occur with an additional power of a Yukawa coupling.

As a result, most of these operators are irrelevant to the decay of the right-handed neutrinos. For example, $\mathcal{O}_3^{(6)}$, $\mathcal{O}_4^{(6)}$ and $\mathcal{O}_7^{(6)}$ can all be ignored since they are suppressed relative to the other operators by at least an additional factor of the lepton or neutrino Yukawa couplings. After electroweak symmetry breaking, the only effect of $\mathcal{O}_6^{(6)}$ is to correct the overall scale of the neutrino coupling by a small amount. It does not contribute to the right-handed neutrino decay since $m_h > m_R$ is assumed throughout. Operator $\mathcal{O}_7^{(6)}$ contributes to the decay of a Higgs boson, but not to the decay of a right-handed neutrino. However, the contribution of this operator to h decay is suppressed by two powers of the neutrino Yukawa coupling in the amplitude and therefore irrelevant.

Of the remaining three operators relevant to the decay of a right-handed neutrino, $\mathcal{O}_1^{(6)}$, $\mathcal{O}_2^{(6)}$ and $\mathcal{O}_5^{(6)}$ contribute to $N \to lW^*$ and $N \to \nu_L Z^*$ decays at $O(\lambda_\nu)$ in the amplitude. This is parametrically the same order as decays caused by neutral and charged current interactions. But since the dimension 6 operators are suppressed by Λ^2, contributions of $\mathcal{O}_1^{(6)}$, $\mathcal{O}_2^{(6)}$ $\mathcal{O}_6^{(6)}$ to these decay processes can be neglected.

However, both $\mathcal{O}_1^{(6)}$ and $\mathcal{O}_2^{(6)}$ contribute to a new decay process. After electroweak symmetry breaking these two operators generate the magnetic moment operator

$$e c_M \frac{1}{\Lambda^2} (N m_D \sigma^{\mu\nu} \nu_L) F_{\mu\nu} \tag{A.17}$$

where $F_{\mu\nu}$ is the electromagnetic field strength and c_M is a linear combination of c_1 and c_2. This operator introduces a new decay mode

$$N \to \gamma \nu_L \ , \ \gamma \overline{\nu}_L \ . \tag{A.18}$$

This decay is further discussed in Sections 2.2.2 and 3.

B. Amplitudes and Decay Rates

Through mass-mixing with the left-handed neutrinos ν_I, the right-handed neutrinos N_I acquires the following couplings to massive gauge bosons and leptons,

$$\mathcal{L}_N = i \frac{g}{\sqrt{2}} D_{JI} W_\mu^- \bar{l}_J \bar{\sigma}^\mu N_I + i \frac{g g_L^{(\nu)}}{\cos\theta_W} D_{JI} g^{(i)} Z_\mu \overline{\nu}_J \bar{\sigma}^\mu N_I + h.c. \tag{B.1}$$

where the "left-handed" and "right-handed" fermion couplings to the Z are

$$\begin{aligned} g_L^{(i)} &= T_3^{(i)} - Q^{(i)} \sin^2\theta_W \\ g_R^{(i)} &= -Q^{(i)} \sin^2\theta_W \end{aligned} \tag{B.2}$$

and where $\sin^2\theta_W \simeq 0.2312$ is the Weinberg angle, $T_3^{(i)} = \pm 1/2$ is the weak isospin , $Q^{(l)} = -1$, $D_{JI} = [m_D^T]_{JI} M_I^{-1}$ is the left-right mixing angle, and two-component notation is used.

For the charged-current mediated decay $N_I \to l_J q\overline{q}'$ one obtains the amplitude

$$A_{CC} = -i\left(\frac{g}{\sqrt{2}}\right)^2 \left[\overline{u}(p_l)\overline{\sigma}^\mu \mu(p_N)\right]\left[\overline{u}(p_q)\overline{\sigma}_\mu \nu(p_{q'})\right] \frac{D_{JI}}{(p_N-p_l)^2 - m_W^2 + i\Gamma_W m_W} \tag{B.3}$$

The spin-averaged matrix element is

$$\frac{1}{2}\sum_{\text{spin}} |\mathcal{A}_{CC}|^2 = |D_{JI}|^2 \frac{32N_c}{v^4} \frac{1}{(1-\frac{M_I^2}{m_W^2}+2\frac{E_l M_I}{m_W^2})^2 + \frac{\Gamma_W^2}{m_W^2}} (p_N\cdot p_q)(p_l\cdot p_{q'}) \tag{B.4}$$

The integral over phase space is the same as in muon decay and be done following [29]. This leads to the exclusive decay rate

$$\Gamma[N_I \to l_J q\overline{q}'] = \frac{G_F^2 M_I^5}{192\pi^3}|D_{JI}|^2 N_c c_W \tag{B.5}$$

where the factor c_W is due to the final integration over the lepton energy and is given by

$$c_G(x_G, y_G) = 2\int_0^1 dz z^2(3-2z)\left((1-(1-z)x_G)^2 + y_G\right)^{-1} \tag{B.6}$$

where $x_G = M_I^2/m_G^2$, $y_G = \Gamma_G^2/m_G^2$, $c_G(0,0) = 1/(1+y_G) \simeq 1$. The non-vanishing momentum transfer enhances the decay rate by approximately 10% for m_R masses around 30GeV, by approximately 50% for masses around 50 GeV. By including the effect of the finite width of the gauge boson this formula can also be used in the region $M_I > m_W$ where the W boson is on-shell.

From the result above one readily obtains

$$\Gamma[N_I \to l_J \bar{l}_{K\neq J}\nu_K] = \frac{G_F^2 M_I^5}{192\pi^3}|D_{JI}|^2 c_W \tag{B.7}$$

the only difference being the absence of a color factor. Decays into same flavor leptons, $K = J$, are dealt with below since there is quantum interference between the charged and neutral currents which must be computed more carefully.

Decays of N_I to $\nu_J q_K \overline{q}_K$, $\nu_J l_{K\neq J}\bar{l}_{K\neq J}$ or to $\nu_J\nu_{K\neq J}\overline{\nu}_{K\neq J}$ are similarly obtained. The only difference is to replace $N_c c_W \to \left[(g_L^{(K)})^2 + (g_R^{(K)})^2\right] N_c^{(K)} c_Z$. Then one finds

$$\Gamma[N_I \to \nu_J f_{K\neq J}\overline{f}_{K\neq J}] = \frac{G_F^2 M_I^5}{192\pi^3}|D_{JI}|^2 \left[(g_L^{(K)})^2 + (g_R^{(K)})^2\right] N_c^{(K)} c_Z \tag{B.8}$$

The amplitude for $N_I \to l_J^- l_J^+ \nu_J$ has interfering contributions from charged and neutral current interactions. The computation of the amplitude and tricky factors of (-1) etc., are similar to the amplitudes for the process $\overline{\nu}e \to \overline{\nu}e$ which is described in, for example [30]. In what follows I parallel their discussion. The amplitude for the charged current contribution is

$$\begin{aligned}\mathcal{A}_{CC} &= -iD_{JI}\left(\frac{g}{\sqrt{2}}\right)^2 \left[\overline{u}(p_{l^-})\overline{\sigma}^\mu u(p_N)\right]\left[\overline{u}(p_\nu)\overline{\sigma}_\mu \nu(p_{l^+})\right]\frac{1}{(p_N-p_{l^-})^2 - m_W^2} \\ &= +iD_{JI}\left(\frac{g}{\sqrt{2}}\right)^2 \left[\overline{u}(p_{l^-})\overline{\sigma}^\mu \nu(p_{l^+})\right]\left[\overline{u}(p_\nu)\overline{\sigma}_\mu u(p_N)\right]\frac{1}{(p_N-p_{l^-})^2 - m_W^2}\end{aligned} \tag{B.9}$$

where a Fierz transformation was performed in the last step. Two-component notation is used here. In ignoring the lepton mass, the coupling of the Z to the left-handed and right-handed lepton currents do not interfere in the overall rate, so I can consider them separately. Only the contribution of the left-handed lepton neutral current to the amplitude interferes with the charged current amplitude. One finds

$$\mathcal{A}_{NC}^{(L)} = iD_{JI}\frac{g^2}{2}\frac{g_L^{(l)}}{\cos^2\theta_W}\left[\overline{u}(p_{l^-})\overline{\sigma}^\mu\nu(p_{l^+})\right]\left[\overline{u}(p_\nu)\overline{\sigma}_\mu u(p_N)\right]\frac{1}{(p_N-p_\nu)^2-m_Z^2} \qquad \text{(B.10)}$$

In the first line of the charged current amplitude there is an overall $(-)$ sign relative to the neutral current amplitude $\mathcal{A}_{NC}^{(L)}$ because of a relative ordering of the anti-commuting creation and annihilation operators that occurs when computing the amplitude. The total is

$$\begin{aligned}\mathcal{A}_{CC}+\mathcal{A}_{NC}^{(L)} &= -i\frac{g^2}{2}D_{JI}\left[\overline{u}(p_{l^-})\overline{\sigma}^\mu\nu(p_{l^+})\right]\left[\overline{u}(p_\nu)\overline{\sigma}_\mu u(p_N)\right] \\ &\quad\times\left[\frac{1}{(p_N-p_l^-)^2-m_W^2}+\frac{g_L^{(l)}}{\cos^2\theta_W}\frac{1}{(p_N-p_\nu)^2-m_Z^2}\right]\end{aligned} \qquad \text{(B.11)}$$

In the low-energy limit the expression in parentheses reduces to

$$\frac{-1}{m_W^2}\left(\frac{1}{2}+\sin^2\theta_W\right) \qquad \text{(B.12)}$$

In this limit the rate is proportional to

$$\left(\frac{1}{2}+\sin^2\theta_W\right)^2 \simeq 0.53\ . \qquad \text{(B.13)}$$

To this must be added the rate for neutral current decays into $l_R^- l_L^+$, which is non-interfering and proportional to $(\sin^2\theta_W)^2 \simeq 0.05$, to obtain a total rate for the $l_J^- l_J^+ \nu_J$ channel which is proportional to 0.59 in the low energy limit. The result (B.13) is somewhat smaller than that obtained by incorrectly incoherently adding the charged and neutral current amplitudes for decays into left-handed charged leptons, which gives $1+(-1/2+\sin^2\theta_W)^2 \simeq 1.07$.

The charged-conjugated decay $N_I \to l_J^- l_J^+ \overline{\nu}_J$ has an identical rate to (B.13). It adds incoherently to the amplitudes above because it has a distinguishable final state; one final state has a neutrino and the other an anti-neutrino.

C. References

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